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(October 9, 2001)

If the underlying flavor symmetry is Abelian, quark mixings in d_R sector are the most prominent. Such flavor violating effects can reveal itself through $\tilde{d}_R$ squark mixings if supersymmetry is realized in Nature. Quark-squark alignment is necessary to deal with Δm_K and ε_K constraints, but interestingly, with $m_{\tilde{q}}$, $m_{\tilde{g}} \sim$ TeV, the $\tilde{d}_R$ mixing effects are comparable to B_d and B_s mixings in the Standard Model, while D^0 mixing is tantalizingly close to some hints from data. CP phases in these mixings would therefore be deviant, and $|V_{td}|$ and $\arg V_{ub}^*$ may be larger than allowed by unitarity constraints, which can be checked by the BaBar and Belle experiments. Mixing induced CP violation in $b \to s\gamma$ and $d\gamma$ transitions can be obtained, in particular, by sizable enhancement with an extra $\tan\beta$ factor from non-standard soft breaking terms. Heavy superparticles can escape present flavor changing neutral current (FCNC) bounds and direct searches at colliders, but reveal themselves in the B system.

PACS numbers: 12.60.Jv, 11.30.Hv, 11.30.Er, 13.25.Hw

I. INTRODUCTION

Despite its great success, the Standard Model (SM) is widely regarded as a weak scale effective theory. We expect to encounter new phenomena that arise from a more complete theory as we increase the luminosity and energy of our probes. Although we have not observed any convincing deviation of experimental results from SM predictions (except neutrino oscillations [1]), there are a few hints on the existence of New Physics.

It is well known that the sine of the CP violating phase, ϕ_1 (or β) $\equiv \arg V_{td}$ (PDG phase convention [2]), can be measured via the time dependent asymmetry,

$$a_{J/\psi K_S^0} = \frac{\Gamma(\bar{B}^0(t) \to J/\psi K_S^0) - \Gamma(B^0(t) \to J/\psi K_S^0)}{\Gamma(\bar{B}^0(t) \to J/\psi K_S^0) + \Gamma(B^0(t) \to J/\psi K_S^0)} = \sin 2\phi_1 \, \sin \Delta m_{B_d} t. \tag{1}$$

Following earlier measurements [3–7], the BaBar and Belle Collaborations have recently firmly established [8,9] $\sin 2\phi_1$ to be nonzero. Compared to earlier low results of $\sin 2\phi_1 = 0.58\,^{+0.32+0.09}_{-0.34-0.10}$ (Belle) [6] and $0.34 \pm 0.20 \pm 0.05$ (BaBar) [7], the corresponding numbers are now $0.99 \pm 0.14 \pm 0.06$ (Belle) [9] and $0.59 \pm 0.14 \pm 0.05$ (BaBar) [8], respectively. Combining the two most recent values without systematic errors, one gets the average value

$$\sin 2\phi_1 = 0.79 \pm 0.10. \tag{2}$$

While this is still consistent with the Cabibbo–Kobayashi–Maskawa unitarity (CKM) fit value of $\sin 2\phi_1 = 0.698 \pm 0.066$ [10], the central value is now somewhat on the high side, especially for the Belle number. If this trend persists — which we would know by summer 2002 — it would imply the presence of New Physics. With this in mind, it is clearly a good time to study other related CP violation processes, especially those that are suppressed in the SM.

It has recently been shown that charmless rare B decays favor [11] a value for ϕ_3 (or γ) $\equiv \arg V_{ub}^*$ that is larger than the one obtained from the CKM unitarity fit [10]. The latter is dominated by recent improved bounds on $\Delta m_{B_s}/\Delta m_{B_d}$ [12],

$$\Delta m_{B_d} = 0.484 \pm 0.010\,\mathrm{ps}^{-1}, \tag{3}$$

$$\Delta m_{B_s} > 15.0\,\mathrm{ps}^{-1} \text{ at 95\% C.L.}, \tag{4}$$

which tends to squeeze out the large ϕ_3 possibility. One also has $\mathrm{Br}(K^+ \to \pi\nu\bar{\nu}) = 4.2^{+9.7}_{-3.5} \times 10^{-10}$ from the E787 Collaboration [13], where the central value is several times above the SM expectation, implying a rather large $|V_{td}|$. The last branching ratio, of course can be viewed as decreasing with time since no new events have been found. Combining the above, however, perhaps a more consistent picture would be if B_d or B_s mixings have additional New Physics sources. This may already be indicated by the measurement of $\sin 2\phi_1$ in Eq. (2) as we have discussed. It can alternatively be tested in the CP phase of B_s mixing, which can be studied at the Tevatron collider in the next few years. If a non-vanishing value is found, it would definitely imply New Physics since the SM prediction is zero.

The E791, CLEO, FOCUS, Belle and BaBar Collaborations have reported search results for CP asymmetries in the neutral D-meson system [14–19]. The search for D^0–$\bar{D}^0$ mixing using CLEO II.V data suggests that $x_D \equiv \Delta m_D/\Gamma$ is less then 2.9% at 95% C.L. [15], which is far below previous results [20]. The actual numbers, however, are $x_D'^2/2 < 0.041\%$ and $-5.8\% < y_D' < 1.0\%$, in terms of

$$x_D' = x_D \cos\delta_D + y_D \sin\delta_D$$
$$y_D' = y_D \cos\delta_D - x_D \sin\delta_D, \tag{5}$$

where δ_D is the relative strong phase between the doubly Cabibbo suppressed $D^0 \to K^+\pi^-$ and the Cabibbo favored $\bar{D}^0 \to K^+\pi^-$ decay amplitudes. The SM predictions of these mixing parameters are small, $x_D \sim 10^{-5} - 10^{-4}$ and $y_D \sim 10^{-4} - 10^{-2}$ [21] (a recent discussion suggests $y_D \sim 1\%$ [22]). The CLEO Collaboration arrived at $|x_D| < 2.9\%$ by assuming δ_D is as small as suggested by SU(3) and other arguments [23,24]. If one removes [25] the prejudice that $D^0 \to K^+\pi^-$ and $K^-\pi^+$ amplitudes have the same strong phase, the result on y'_D may actually be hinting at a sizable x_D, which would strongly suggest short distance New Physics interactions. It is therefore important to compare x'_D, y'_D with other D^0–$\bar{D}^0$ mixing related measurements.

The CP asymmetry parameter y_{CP} is related to the lifetime difference between $D^0 \to K^-\pi^+$ and $D^0 \to K^-K^+$, where the former is flavor specific and the latter a CP eigenstate. The FOCUS Collaboration had found the intriguing value of $y_{CP} = (3.42 \pm 1.39 \pm 0.74)\%$ [16], which has a significance of more than 2σ. However, recently, the CLEO, Belle and BaBar Collaborations find much lower values. The current world average is $(1.1 \pm 0.9)\%$ [22]. Still, the point remains that one could have large x_D if $\sin\delta_D$ is sizable [26].

The processes $b \to s\gamma, d\gamma$ occur only at loop level in the SM, therefore they are naturally sensitive to New Physics. The processes $B \to K^*\gamma$ and $b \to s\gamma$ have long been observed [27]. However, the Belle Collaboration observes a 3.5 ± 2.1 event excess for $B \to \rho\gamma$, giving $\mathrm{Br}(B \to \rho^0\gamma) < 10.6 \times 10^{-6}$ [28] at 90% C.L. The quark level $bq\gamma$ coupling is usually parametrized as

$$H_{\rm SM} = -\frac{G_F}{\sqrt{2}}\frac{e}{4\pi^2} V_{tb}V_{tq}^* \, m_b \, \bar{q}[C_{7\gamma}R + C'_{7\gamma}L]\sigma_{\mu\nu}F^{\mu\nu}b + h.c., \qquad (6)$$

where in the SM, $C_{7\gamma} \cong -0.3$ at the typical B decay energy scale $\mu \approx 5$ GeV. Due to the left-handed nature of weak interactions, $C_{7\gamma}$ dominates while $C'_{7\gamma}$ is suppressed by m_q/m_b. This may not be the case, however, in models beyond the SM, and interesting CP violating asymmetries in mixing induced radiative B decay can occur [29]. In a previous work [30], we have discussed the $b \to s\gamma$ process in the context of New Physics. We found that large direct CP violations are possible, in contrast to SM expectations which are very small. For mixing induced CP violation in B decay, we also find large asymmetries due to enhanced $C'_{7\gamma}$. Furthermore, the mixing dependent asymmetry in $B \to \rho\gamma$ is more accessible than in $B \to K^{*0}\gamma$, because $\rho^0 \to \pi^+\pi^-$ can give the vertex information needed for time dependence, while $K^{*0} \to K_s\pi^0$ unfortunately does not [29]. Given that SM predictions on direct CP violation in $b \to d\gamma$ are in general not small [31], mixing induced CP violation in $b \to d\gamma$ is a more sensitive probe of New Physics. Of course, mixing dependent CP violation in $b \to s\gamma$ processes can be more readily probed in B_s system such as in $B_s \to \phi\gamma$.

Mixing induced CP violation in radiative B decay is similar to the golden $J/\psi K_S$ mode. The hadronic uncertainties factor out for self-conjugate CP eigenstate(s). Furthermore, one needs both B and $\bar{B}$ to decay to the same final state to allow interference to take place. Therefore, the formula of the asymmetry $a_{M^0\gamma}$, where M^0 is a CP self-conjugate hadron such as ρ^0, ω, ϕ, K^{*0}(in $K_s\pi^0$), resembles that of $a_{J/\psi K_s^0}$,

$$a_{M^0\gamma} = \xi \sin 2\vartheta \, \sin[2\phi_1 - \phi_{7\gamma} - \phi_{7\gamma}^{(\prime)}] \, \sin\Delta m \, t \qquad (7)$$

where ξ is the CP eigenvalue of M^0,

$$\sin 2\vartheta \equiv \frac{2\,|C_{7\gamma}C'_{7\gamma}|}{|C_{7\gamma}|^2 + |C'_{7\gamma}|^2}, \qquad (8)$$

is the mean strength of the two chiral amplitudes, and $\phi_{7\gamma}^{(\prime)}$ is the phase of $C_{7\gamma}^{(\prime)}$. While the "$\sin 2\vartheta$" in the $B \to J/\psi K_S$ case is equal to one, an important feature for $B \to M^0\gamma$ is that one needs both $C_{7\gamma}$ and $C'_{7\gamma}$ for the interference to occur. This is the reason why these asymmetries are suppressed in the SM since $|C'^{\rm SM}_{7\gamma}/C^{\rm SM}_{7\gamma}| \ll 1$, hence precisely why they are sensitive to New Physics. The Belle Collaboration may be able to test the asymmetry to an accuracy of 10% in about a decade [32]. Hadron machines may also be able to study this if they can observe the radiative B decay modes, since they produce many more $B\bar{B}$s than the $e^+\,e^-$ machines.

Another hint for New Physics may come from the rather large value of the newly observed ε'/ε [33]. It has prompted New Physics considerations [34], even though the experimental value could be accommodated within the SM. We do not consider New Physics hints from muon anomalous magnetic moment.

We are interested in New Physics models that can lead to deviations in the above mentioned processes. It is interesting that supersymmetry (SUSY) models with Abelian horizontal symmetry (AHS) can provide a unified framework for all such New Physics effects. This will be presented in Section II, where we will explain the implications of AHS on quark mixing, in particular on the origin of large right-handed down quark mixings. The effect is carried over to the squark sector in SUSY models, and squark mixings will impact on flavor changing neutral currents (FCNC). In face of stringent constraints from kaon mixings, quark-squark alignment (QSA) is invoked to produce texture zeros, where we will give an explicit example of horizontal charge assignments. The subsequent sections are devoted to various FCNCs induced by squark mixings. In Section III we study the effect on B_d–$\bar{B}_d$ and B_s–$\bar{B}_s$ mixings, which sets the scale of superparticle masses. In Section IV, we show that the chargino contribution on kaon mixings gives a similar superparticle mass scale. A generic feature of QSA is the possibility of sizable D^0–$\bar{D}^0$ mixing. It is interesting that in AHS models with SUSY, with sparticle scale fixed by B–$\bar{B}$ mixing, x_D could be right in the ball-park of the CLEO range, as we show in Section V. Section VI is devoted to radiative B decay in SUSY models, and discussions and conclusions are given in Sections VII and VIII, respectively.

II. ABELIAN FLAVOR SYMMETRY WITH SUSY

A. Abelian Flavor Symmetry and Large $\tilde{d}_R$ Mixings

Fermion masses and mixings in the Cabibbo-Kobayashi-Maskawa matrix $V_{\rm CKM}$ exhibit an intriguing hierarchical pattern in powers of $\lambda \equiv |V_{us}|$:

$$\begin{aligned} m_u/m_c \sim \lambda^3, \quad m_c/m_t \sim \lambda^4, \quad m_t\sqrt{G_F} \sim 1, \\ m_d/m_s \sim \lambda^2, \quad m_s/m_b \sim \lambda^2, \quad m_b/m_t \sim \lambda^3, \\ |V_{cb}| \sim \lambda^2, \qquad |V_{ub}| \sim \lambda^3. \end{aligned} \tag{9}$$

The structure could be due to an underlying symmetry [35], the breaking of which gives an expansion in $\lambda \sim \langle S\rangle/M$, with S a scalar field and M a high scale. In these models, after integrating out some massive fields of mass M, one obtains non-renormalizable terms [35,36],

$$L_{\rm mass} = \lambda^q_{ij} H_d \left(\frac{S}{M}\right)^{\alpha_{q\,ij}} Q_i \bar{q}_{Rj} + {\rm h.c.}, \tag{10}$$

where q are summed over up and down type quarks, λ^q are $\mathcal{O}(1)$ numbers, and i, j are generation indices. $L_{\rm mass}$ is made a horizontal symmetry singlet by choosing appropiate powers of S, i.e. $\alpha_{q\,ij}$. For models with Abelian horizontal symmetry, without loss of generality [36], we can define the horizontal charges of the scalar fields as

$$H(H_d) = H(H_u) = 0,\; H(S) = -1. \tag{11}$$

The breaking of the horizontal symmetry as well as electroweak symmetry lead to quark mass elements,

$$M^*_{q\,ij} = \lambda^q_{ij} \langle H_q\rangle \left(\frac{\langle S\rangle}{M}\right)^{\alpha_{q\,ij}}, \;\alpha_{q\,ij} = H(Q_i) + H(\bar{q}_{Ri}). \tag{12}$$

It is now easy to see that

$$M_{ij}M_{ji} \sim M_{ii}M_{jj},\; (i,\, j \text{ not summed}), \tag{13}$$

which follows as a consequence of the commuting nature of horizontal charges.

By assuming small mixing angles, as one tries to understand the hierarchical pattern in λ by AHS, the quark mass ratios fix the order of magnitude of the diagonal elements of quark mass matrices. The upper right part of the mass matrix M_q corresponds to U_{qL} rotation, which is related to $V_{\rm CKM} = U_{uL}U^\dagger_{dL}$. Small mixing angle and naturalness imply $U_{qL} \sim V_{\rm CKM}$. By using Eq. (13), one can work out the lower left part and hence the whole mass matrix [36,37]

$$\frac{M_u}{m_t} \sim \begin{bmatrix} \lambda^7 & \lambda^5 & \lambda^3 \\ \lambda^6 & \lambda^4 & \lambda^2 \\ \lambda^4 & \lambda^2 & 1 \end{bmatrix}, \quad \frac{M_d}{m_b} \sim \begin{bmatrix} \lambda^4 & \lambda^3 & \lambda^3 \\ \lambda^3 & \lambda^2 & \lambda^2 \\ \lambda & 1 & 1 \end{bmatrix}. \tag{14}$$

Since U_{qL} are restricted by $V_{\rm CKM}$, mixing angles in U_{qR} are in general greater. In particular, we find that M^{32}_d/m_b and M^{31}_d/m_b are the most prominent off-diagonal elements.

To summerize, we have,

$$U_{qL} \sim \begin{pmatrix} 1 & \lambda & \lambda^3 \\ \lambda & 1 & \lambda^2 \\ \lambda^3 & \lambda^2 & 1 \end{pmatrix}, \quad U_{dR} \sim \begin{pmatrix} 1 & \lambda & \lambda \\ \lambda & 1 & 1 \\ \lambda & 1 & 1 \end{pmatrix},$$
$$U_{uR} \sim \begin{pmatrix} 1 & \lambda & \lambda^4 \\ \lambda & 1 & \lambda^2 \\ \lambda^4 & \lambda^2 & 1 \end{pmatrix}. \tag{15}$$

It is clear that mixing angles in U_{dR} are in general greater than those in U_{qL} and U_{uR}, especially when the b flavor is involved. * Although large mixings in the right handed down quark sector are *useless* or well hidden within SM, B_d and B_s mixings are naturally susceptible to New Physics involving further dynamics related to the right-handed down flavor sector.

As one of the leading candidates for New Physics, SUSY helps resolve many of the potential problems that emerge when one extends beyond the SM, for example the gauge hierarchy problem, unification of SU(3)×SU(2)×U(1) gauge couplings, and so on [39]. It is interesting to note that large mixing in right-handed down quark sector will be transmitted to right-handed down squarks, if the breakings of flavor symmetry and SUSY are not closely related. The flavor symmetry gives better control on soft breaking parameters resulting in a more predictive SUSY model. We now elevate Eq. (10) to the superpotential as well as similar forms in trilinear terms, the so called A terms. By flavor symmetry, we still have the same power of S to ensure that the whole term remains a horizontal singlet, that is,

$$(\widetilde{M}^2_q)^{ij}_{LR} \sim \widetilde{m}\, M^{ij}_d, \quad (\widetilde{M}^2_q)_{RL} = (\widetilde{M}^2_q)^\dagger_{LR}, \tag{16}$$

which are *roughly* proportional to respective quark mass matrices, hence their effects are suppressed by $m_q/\widetilde{m}$ [37]. While the symmetry does not require the new λ_{ij} in the A-term to be the same as λ^q_{ij} in Eq. (10), $(\widetilde{M}^2_q)_{LR}$ cannot, in general, be diagonal in the quark mass basis. From Eq. (14), one easily gets $(\widetilde{M}^2_Q)_{LL}/\widetilde{m}^2 \sim V_{\rm CKM}$, while

$$(\widetilde{M}^2_d)_{RR} \sim \widetilde{m}^2 \begin{bmatrix} 1 & \lambda & \lambda \\ \lambda & 1 & 1 \\ \lambda & 1 & 1 \end{bmatrix}. \tag{17}$$

The squark mixings will have impact on FCNC because of extra dynamics involving $\tilde{q}$-$\tilde{g}$, $\tilde{q}$-$\tilde{\chi}^\pm$ and $\tilde{q}$-$\tilde{\chi}^0$ couplings. It is now clear that FCNC processes involving $\tilde{b}_R$ are sensitive probes of this generic class of SUSY models.

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*These large mixings in second and third generation d_R's may be related to large mixing in second and third generation neutrinos [1] when considering grand unified theories (GUT) [38].

The RR sector could contribute significantly to B_d and B_s mixings, to be discussed in the next section, via large mixings in $\tilde{b}_R$–$\tilde{d}_R$ and in $\tilde{b}_R$–$\tilde{s}_R$ as shown in Eq. (17).

B. Quark-squark Alignment

In order to compare squark mixing angles with FCNC constrains, we will use the mass insertion approximation [40,41] in the following discussion. It is customary to take squarks as almost degenerate at scale $\widetilde{m}$. In quark mass basis, one defines [40],

$$\delta^{ij}_{qAB} \equiv [U^\dagger_{qA}\,(\widetilde{M}^2_q)_{AB}\,U_{qB}]^{ij}/\widetilde{m}^2, \tag{18}$$

which is roughly the squark mixing angle, $\widetilde{M}^2_q$ are squark mass matrices, $A, B = L, R$, and i, j are generation indices. Note that $\delta^{13}_{dRR} \sim \lambda$ and $\delta^{23}_{dRR} \sim 1$, while LR and RL mixings are suppressed by $m_q/\widetilde{m}$.

It is well known that kaon mixings give stringent constraint on new flavor violating source. The $12, 21$ elements in $\widetilde{M}^2_Q$ and $(\widetilde{M}^2_d)_{RR}$ in Eq. (17) will induce too large a contribution to kaon mixing via gluino box diagrams [42,43]. One has to suppress these squark mixings by enforcing approximate "texture zeros". This can be done by invoking quark-squark alignment (QSA) [36,37], by using two (or more) singlet fields S_i to break the $U(1) \times U(1)$ (or higher) Abelian horizontal symmetry, and making use of the holomorphic nature of the superpotential in SUSY models. For example, we may have a term with negative power α_{qij} in S in Eq. (10) to satisfy the horizontal symmetry. The term $S^{-|\alpha_{qij}|}$ is simply $(S^*)^{|\alpha_{qij}|}$. As we promote Eq. (10) to superpotential, which can only be a function of superfields and not of conjugate superfields at the same time, one can no longer use S^* and hence there is a zero in that particular ij-th element. One can have $M^{12,21}_d = 0$ which imply $U^{12}_{dL,R} = 0$ or are highly suppressed, and likewise $(\widetilde{M}^2_d)^{12}_{LL,RR}$ are also suppressed. Thus, $\delta^{12}_{dLL,RR}$ can be suppressed and the kaon mixing constraint is satisfied accordingly.

There is one subtlety involving our choice to retain $(\widetilde{M}^2_d)^{13}_{RR}$, which arises from M^{31}_d. The mass matrix M_d is diagonalized by a bi-unitary transform, hence, it is of the form

$$\frac{M_d}{m_b} = U^\dagger_{dL}\,\frac{M^{\rm diag}_d}{m_b}\,U_{dR}$$
$$\sim \begin{pmatrix} 1 & \lambda^a & \lambda^b \\ \lambda^a & 1 & \lambda^c \\ \lambda^b & \lambda^c & 1 \end{pmatrix} \begin{pmatrix} \lambda^4 & 0 & 0 \\ 0 & \lambda^2 & 0 \\ 0 & 0 & 1 \end{pmatrix} \begin{pmatrix} 1 & \lambda^d & \lambda^e \\ \lambda^d & 1 & \lambda^f \\ \lambda^e & \lambda^f & 1 \end{pmatrix} \tag{19}$$

where the diagonal matrix in the middle of the right hand side corresponds to the diagonal down quark mass (ratio) matrix, and the one to its left (right) corresponds to U_{dL} (U_{dR}). Multiplying out the matrices in Eq. (19), we have

$$\frac{M^{11}_d}{m_b} \sim \lambda^4 + \lambda^{2+a+d} + \lambda^{b+e},\ \frac{M^{12}_d}{m_b} \sim \lambda^{4+d} + \lambda^{2+a} + \lambda^{b+f},$$
$$\frac{M^{21}_d}{m_b} \sim \lambda^{4+a} + \lambda^{2+d} + \lambda^{c+e},\ \frac{M^{22}_d}{m_b} \sim \lambda^{4+a+d} + \lambda^2 + \lambda^{c+f},$$
$$\frac{M^{31}_d}{m_b} \sim \lambda^{4+b} + \lambda^{2+c+d} + \lambda^e,\ \frac{M^{32}_d}{m_b} \sim \lambda^{4+b+d} + \lambda^{2+c} + \lambda^f,$$
$$\frac{M^{13}_d}{m_b} \sim \lambda^{4+e} + \lambda^{2+a+f} + \lambda^b,\ \frac{M^{23}_d}{m_b} \sim \lambda^{4+a+e} + \lambda^{2+f} + \lambda^c,$$
$$\frac{M^{33}_d}{m_b} \sim 1 + \lambda^{4+b+e} + \lambda^{2+c+f}. \tag{20}$$

We see that, by retaining $M^{31}_d/m_b \sim \lambda$, we have $e = 1$ and hence $U^{13}_{dR} \sim \lambda$. We will have $c = 2$, if we also retain $M^{23}_d/m_b \sim \lambda^2$. However, as the kaon mixing constraint requires

$$M^{21}_d/m_b \sim \lambda^{4+a} + \lambda^{2+d} + \lambda^{c+e} \sim 0, \tag{21}$$

which in turn gives $d = 1$ for λ^{2+d} to be of the same order as $\lambda^{c+e} = \lambda^3$, to allow cancellation to take place. This will still give squark mixing angle $\sim \lambda$ between $\tilde{d}_R$–$\tilde{s}_R$, which is not acceptable. A closer look reveals that $M^{32}_d/m_b \sim 1$ is also unacceptable. It gives $f = 0$ from $M^{32}/m_b \sim 1$. With the requirement of $M^{23}_d = 0$, we again have $c = 2$ and it follows that $d = 1$ by the same argument. Because $(\widetilde{M}^2_d)^{13}_{RR}/\widetilde{m}^2 \sim \lambda$ is kept, we need to make M^{23}_d/m_b *and* M^{32}_d/m_b also vanishing in face of stringent Δm_K and ε_K constraints. The decoupling of s flavor from other generations thus follows from imposing QSA in 12 sector *and* choosing to retain $M^{31}_d \neq 0$.

In the usual approach of quark-squark alignment, this subtlety does not arise because one aspires to lower $m_{\tilde{q}}$, $m_{\tilde{g}}$ for sake of collider and other signatures. Since $\delta^{13}_{dRR} \sim \lambda$ and $\delta^{23}_{dRR} \sim 1$ would then violate B_d mixing and $b \to s\gamma$ constraints already, they are eliminated from the outset. As a result, $M^{23}_d/m_b \sim \lambda^2$, though of little consequence, can be retained.

C. Explicit Examples of QSA

To be specific, we now give an explicit assignment of the horizontal charges of quark superfields and the resulting mass matrices as an illustrative example.

Since $|V_{ub}| \sim 0.002 - 0.005 < \lambda^3$ [44], we may use a smaller parameter $\tilde{\lambda} = 0.18$ instead of λ. We use two S_i fields to break the horizontal symmetry,

$$\frac{\langle S_1\rangle}{M} \sim \tilde{\lambda}^{0.5}, \quad \frac{\langle S_2\rangle}{M} \sim \tilde{\lambda}^{0.5}. \tag{22}$$

The horizontal charges of S_1 and S_2 are $(-1, 0)$, $(0, -1)$, respectively and those of Q, $\bar{d}_R$ and $\bar{u}_R$ are given by

$$\begin{aligned} &Q_1 : (8, -2), \quad Q_2 : (1, 3), \quad Q_3 : (2, -2), \\ &\bar{d}_{R1} : (-2, 10), \quad \bar{d}_{R2} : (9, -3), \quad \bar{d}_{R3} : (-2, 8), \\ &\bar{u}_{R1} : (-3, 11), \quad \bar{u}_{R2} : (0, 3), \quad \bar{u}_{R3} : (-1, 2), \end{aligned} \tag{23}$$

(q,ij)	$\lvert\delta^{ij}_{qLL}\rvert$	$\lvert\delta^{ij}_{qRR}\rvert$	$\lvert\delta^{ij}_{qLR}\rvert$	$\lvert\delta^{ij}_{qRL}\rvert$
(d,12)	$\tilde\lambda^6$	$\tilde\lambda^{12[7]}$	$m_b/\widetilde m\,\tilde\lambda^8(\underline{1}+\{\tan\beta\})$	$m_b/\widetilde m\,(\underline{\tilde\lambda^{14[7]}}+\{\tilde\lambda^4\tan\beta\})$
(d,13)	$\tilde\lambda^3$	$\tilde\lambda$	$m_b/\widetilde m\,\tilde\lambda^3(1+\{\tan\beta\})$	$m_b/\widetilde m\,\tilde\lambda\,(1+\{\tan\beta\})$
(d,23)	$\tilde\lambda^3$	$\tilde\lambda^{11[6]}$	$m_b/\widetilde m\,\tilde\lambda^3(\underline{1}+\{\tan\beta\})$	$m_b/\widetilde m\,(\underline{\tilde\lambda^{11[7]}}+\{\tilde\lambda^5\tan\beta\})$
(u,12)	$\tilde\lambda$	$\tilde\lambda^{4.5}$	$m_t/\widetilde m\,\tilde\lambda^4$	$m_t/\widetilde m\,(\underline{\tilde\lambda^{8.5}}+\{\tilde\lambda^{7.5}\})$
(u,23)	$\tilde\lambda^2$	$\tilde\lambda$	$m_t/\widetilde m\,\tilde\lambda^2$	$m_t/\widetilde m\,\tilde\lambda$

TABLE I. The order of magnitudes of $|\delta^{ij}_q|$s from the Abelian horizontal symmetry model, where $\tilde\lambda = 0.18$. Terms in parentheses [...] correspond to $\tan\beta \sim 50$ case and terms with {...} only exist when we consider non-standard soft breaking terms. The underlined terms are from filled zeros (see text).

for small $\tan\beta$. For $\tan\beta \sim 50$ we change horizontal charges of $\bar d_{Ri}$ to,

$$\bar d_{R1} : (-2,5), \quad \bar d_{R2} : (4,-3), \quad \bar d_{R3} : (-2,3). \tag{24}$$

In this way we get,

$$\frac{M_u}{m_t} \sim \begin{pmatrix} \tilde\lambda^{6.5} & \tilde\lambda^4 & \tilde\lambda^3 \\ 0 & \tilde\lambda^3 & \tilde\lambda^2 \\ 0 & \tilde\lambda & 1 \end{pmatrix}, \quad \frac{M_d}{m_t\,\tilde\lambda^{2.5}} \sim \begin{pmatrix} \tilde\lambda^4 & 0 & \tilde\lambda^3 \\ 0 & \tilde\lambda^2 & 0 \\ \tilde\lambda & 0 & 1 \end{pmatrix}. \tag{25}$$

From Eq. (25), we have effectively decoupled the second generation from the first and third in M_d, which corresponds to suppressed $U^{12,23}_{dL,dR} = 0$. The case is reminiscent of [30] where we decouple d flavor. The corresponding squark mass matrices from Eq. (23) are $(\widetilde M^2_q)^{ij}_{LR} \sim \widetilde m\, M^{ij}_q$ and

$$(\widetilde M^2_Q)^{ij}_{LL} \sim \widetilde m^2 \begin{pmatrix} 1 & \tilde\lambda^6 & \tilde\lambda^3 \\ \tilde\lambda^6 & 1 & \tilde\lambda^3 \\ \tilde\lambda^3 & \tilde\lambda^3 & 1 \end{pmatrix}, \tag{26}$$

$$(\widetilde M^2_u)^{ij}_{RR} \sim \widetilde m^2 \begin{pmatrix} 1 & \tilde\lambda^{4.5} & \tilde\lambda^{4.5} \\ \tilde\lambda^{4.5} & 1 & \tilde\lambda \\ \tilde\lambda^{4.5} & \tilde\lambda & 1 \end{pmatrix}, \tag{27}$$

$$(\widetilde M^2_d)^{ij}_{RR} \sim \widetilde m^2 \begin{pmatrix} 1 & \tilde\lambda^{12} & \tilde\lambda \\ \tilde\lambda^{12} & 1 & \tilde\lambda^{11} \\ \tilde\lambda & \tilde\lambda^{11} & 1 \end{pmatrix}. \tag{28}$$

For large $\tan\beta$, we change $(\widetilde M^2_d)^{12}_{RR}$ and $(\widetilde M^2_d)^{23}_{RR}$ to $\tilde\lambda^7$ and $\tilde\lambda^6$, respectively. We summarize all δ's of interest in Table I. We will see that these values are all well below the limits from Δm_K and ε constraints, even with $\mathcal{O}(1)$ phases.

At this point one thing needs to be emphasized. In face of severe kaon mixing constraints, we did not choose horizontal charges to *create* large right handed $\tilde d$–$\tilde b$ squark mixings. Instead, the choice of horizontal charges were to *retain* this natural large mixing that follow from the Abelian nature of the underlying flavor symmetry. Both the Abelian flavor symmetry and the mixing pattern originate from the observed mass mixing hierarchy pattern. Phenomenological consequences of these large mixings [45] should be explored [46].

There is a generic feature [36,37] of QSA that is worthy of note. Having $U^{12}_{dL} \simeq 0$ implies $U^{12}_{uL} \sim |V_{cd}| = \lambda$, which can also be read off from Eq. (25). One now has $\delta^{12}_{uLL} \sim \tilde\lambda$, as one can see from Eq. (18). This Cabibbo strength δ^{12}_{uLL} can contribute to kaon mixing via chargino diagrams, and also D^0–$\bar D^0$ mixing via gluino diagrams, as we will dicuss in Sec. IV and V, respectively. We note that New Physics contributions to D^0–$\bar D^0$ mixing are of great interest at present, since the recent CLEO (and FOCUS) results may be a hint for D^0 mixing in disguise. Note also that the texture zeros of $M^{21,31}_u$ are generated through QSA. The zero of M^{21}_u is needed to avoid $\delta^{12}_{uRR} \sim \tilde\lambda$, for otherwise, together with $\delta^{12}_{uLL} \sim \tilde\lambda$ they will induce too large a contribution to D mixing. The zero of M^{31}_u is to avoid a feed back to δ^{12}_{uRR}, analogous to the discussion in the previous subsection.

We mention another subtlety arising from the Kähler potential [45]. When the horizontal symmetry is spontaneously broken, mixing also occurs in the kinetic terms. By canonical normalization of the kinetic terms, further mixings are introduced. For example, M_q now becomes $L_q\, M_q\, R^\dagger_q$, where $L_q \sim (\widetilde M^2_Q/\widetilde m^2)^{-1/2}$ and $R^\dagger_q \sim (\widetilde M^2_{qRR}/\widetilde m^2)^{-1/2}$. The zeros in Eq. (25) are now all lifted, and are called filled zeros [45], giving rise to the underlined terms in Table I. U_{qL} also becomes $L_q\, U_{qL}$ and similarly for other mixing matrices. Modifications of previous results can be achieved by suitable rotations and are also shown in Table I.

A second possibility of horizontal charge assignment is to retain M^{32}_d and M^{23}_d while having vanishing M^{31}_d, i.e.

$$\frac{M_d}{m_b} \sim \begin{pmatrix} \lambda^4 & 0 & 0 \\ 0 & \lambda^2 & \lambda^2 \\ 0 & 1 & 1 \end{pmatrix}. \tag{29}$$

The assignment of horizontal charges for this case can be found in Ref. [36]. The squark mixing can generate sizable contribution in B_s mixings [47]. In this case one may have large CP phase in B_s mixing and possible effects in $b \to s\gamma$ to be discussed later. The democratic structure of $\widetilde M^2_{dRR}$ in Eq. (17) in the 2-3 sub-matrix leads to approximate maximal mixing in $\tilde s_R$ and $\tilde b_R$. The large off-diagonal elements $(\widetilde{M}_d^2)_{RR}^{23,32}$ lead to large level splitting. This allows for a possibly light strange beauty squark with interesting impact in B_s mixing and direct search, but leaving $\mathrm{Br}(B \to X_s\gamma)$ largely unaffected [48].

It is interesting that, faced with stringent kaon constraint, AHS models with QSA allow large mixing in either $\tilde{b}_R - \tilde{d}_R$ or $\tilde{b}_R - \tilde{s}_R$, but not both at the same time. Thus, a prediction of this model is, if it is responsible for the smallness of the measured $\sin 2\phi_1$ (assuming that the low value persists in the future), there will be no large New Physics contribution to B_s mixing.

We now study the FCNC induced by these squark mixings in the following sections.

III. B^0–$\bar{B}^0$ MIXING

In this section, we first focus on the general formalism of neutral B meson mixings in the AHS model with SUSY. We will focus on the B_d system for applications, which is readily extendable to the B_s system. For the latter system, the intriguing possibility that large right-handed squark mixings could lead to a light "strange-beauty" squark will be discussed briefly in Sec. VII.C.

The effective Hamiltonian for B_q^0–$\bar{B}_q^0$ mixings from SUSY contributions, where $q = d$ or s, is given by

$$H_{\mathrm{eff}} = \sum_i C_i \mathcal{O}_i, \tag{30}$$

where,

$$\begin{aligned}
&\mathcal{O}_1 = \bar{q}_L^\alpha \gamma_\mu b_L^\alpha \, \bar{q}_L^\beta \gamma^\mu b_L^\beta, \\
&\mathcal{O}_2 = \bar{q}_L^\alpha b_R^\alpha \, \bar{q}_L^\beta b_R^\beta, \quad \mathcal{O}_3 = \bar{q}_L^\alpha b_R^\beta \, \bar{q}_L^\beta b_R^\alpha, \\
&\mathcal{O}_4 = \bar{q}_L^\alpha b_R^\alpha \, \bar{q}_R^\beta b_L^\beta, \quad \mathcal{O}_5 = \bar{q}_L^\alpha b_R^\beta \, \bar{q}_R^\beta b_L^\alpha,
\end{aligned} \tag{31}$$

together with three other operators $\tilde{\mathcal{O}}_{1,2,3}$ that are chiral conjugations ($L \leftrightarrow R$) of $\mathcal{O}_{1,2,3}$. The Wilson coefficients receive charged Higgs, chargino, gluino, gluino-neutralino, and neutralino exchange box diagram contributions,

$$C_i = C_i^{H^-} + C_i^{\tilde{\chi}^-} + C_i^{\tilde{g}} + C_i^{\tilde{g}\tilde{\chi}^0} + C_i^{\tilde{\chi}^0}, \tag{32}$$

where the Feynman diagrams are shown in Fig. 1.

A. Formulas

<u>Charged Higgs box</u> [49,50]:

$$\begin{aligned}
C_1^{H^-} = \frac{\alpha_{\mathrm{w}}^2}{8m_W^2}(V_{tb}V_{tq}^*)^2 \Big[& x_{t\mathrm{W}}\, x_{tH} \cot^4\beta \frac{1}{4} G(x_{tH}, x_{tH}) \\
& + 2x_{t\mathrm{W}}^2 \cot^2\beta \Big(F'(x_{t\mathrm{W}}, x_{t\mathrm{W}}, x_{H\mathrm{W}}) \\
& + \frac{1}{4} G'(x_{t\mathrm{W}}, x_{t\mathrm{W}}, x_{H\mathrm{W}}) \Big) \Big],
\end{aligned} \tag{33}$$

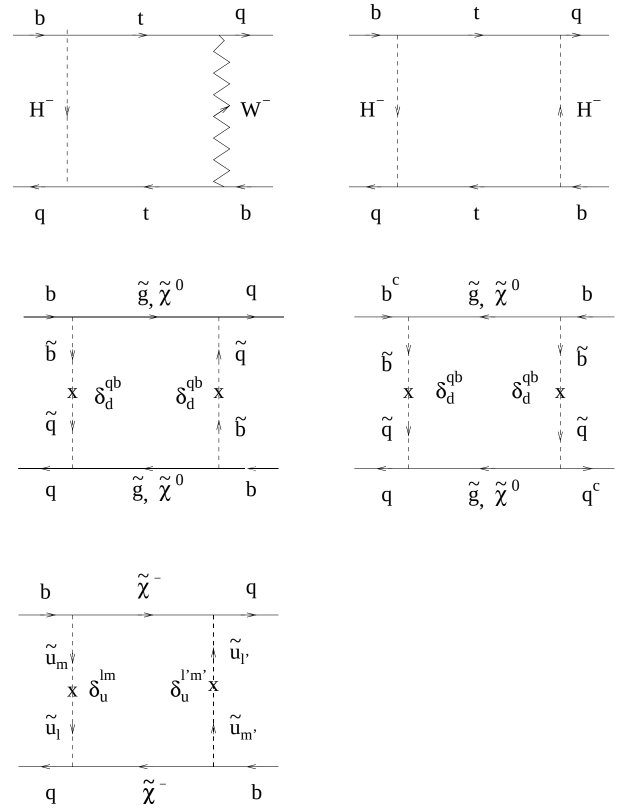

FIG. 1. SUSY box diagrams for $\Delta B = 2$ processes.

$$\begin{aligned}
C_2^{H^-} = -\frac{\alpha_{\mathrm{w}}^2}{8m_W^2}(V_{tb}V_{tq}^*)^2 x_{b\mathrm{W}} \big[& x_{t\mathrm{W}} x_{tH} F(x_{tH}, x_{tH}) \\
& + 2x_{t\mathrm{W}}\, x_{t\mathrm{W}} \cot\beta\, F'(x_{t\mathrm{W}}, x_{t\mathrm{W}}, x_{H\mathrm{W}}) \big],
\end{aligned}$$

where $x_{ab} \equiv m_a^2/m_b^2$, and the loop functions $F^{(\prime)}, G^{(\prime)}$ are given in Ref. [49]. The charge Higgs contributions to other Wilson coefficients are either vanishing or suppressed by m_q/m_W.

<u>Chargino box</u>:

$$\begin{aligned}
C_1^{\tilde{\chi}^-} &= \frac{\alpha_{\mathrm{w}}^2}{8m_{\tilde{\chi}_k^-}^2} \widetilde{G}'(x_{\tilde{q}\tilde{\chi}_k^-}, x_{\tilde{\chi}_j^-\tilde{\chi}_k^-}) \mathcal{A}_{jk}\, \mathcal{A}_{kj} \\
C_3^{\tilde{\chi}^-} &= -\frac{\alpha_{\mathrm{w}}^2}{2m_{\tilde{\chi}_k^-}^2} \sqrt{x_{\tilde{\chi}_j^-\tilde{\chi}_k^-}}\, \widetilde{F}'(x_{\tilde{q}\tilde{\chi}_k^-}, x_{\tilde{\chi}_j^-\tilde{\chi}_k^-}) \\
&\qquad \mathcal{U}_{j2}\, \mathcal{U}_{k2}\, \hat{Y}_b^2 \mathcal{B}_j \mathcal{B}_k
\end{aligned} \tag{34}$$

where the indices j, k are summed over 1 to 2, and

$$\begin{aligned}
\mathcal{A}_{jk} &\equiv \mathcal{V}_{j1}\mathcal{V}_{k1}^* V_{lq}^* V_{mb} \delta_{uLL}^{lm} - \mathcal{V}_{j2}\mathcal{V}_{k1}^* \hat{Y}_t V_{tq}^* V_{tb} \delta_{uRL}^{33} \\
&\quad - \mathcal{V}_{j1}\mathcal{V}_{k2}^* \hat{Y}_t V_{tq}^* V_{tb} \delta_{uLR}^{33} + \mathcal{V}_{j2}\mathcal{V}_{k2}^* \hat{Y}_t^2 V_{lq}^* V_{mb} \delta_{uRR}^{33}, \\
\mathcal{B}_j &\equiv \mathcal{V}_{j1} V_{lq}^* V_{mb} \delta_{uLL}^{lm} - \mathcal{V}_{j2} \hat{Y}_t V_{tq}^* V_{tb} \delta_{uRL}^{33},
\end{aligned} \tag{35}$$

the indices l, m are summed over 3 generations of up type squarks, $\hat{Y}_{u,c,t} = m_{u,c,t}/(\sqrt{2}m_W \sin\beta)$ and similarly, $\hat{Y}_{d,s,b} = m_{d,s,b}/(\sqrt{2}m_W \cos\beta)$. The chargino mixing matrices $\mathcal{U},\mathcal{V}$ in Eq. (34) diagonalize the chargino mass matrix,

$$M_{\widetilde{\chi}^\pm}=\mathcal{U}^*\begin{pmatrix} M_2 & \sqrt{2}m_W\sin\beta \\ \sqrt{2}m_W\cos\beta & \mu\end{pmatrix}\mathcal{V}^\dagger, \qquad (36)$$

and

$$\begin{aligned}&\Big(\widetilde{F}'(x,y),\,\widetilde{G}'(x,y)\Big)\\ &=x^2\partial_a\partial_b\Big(F'(a,b,y),G'(a,b,y)\Big)\Big|_{a=b=x}, \qquad (37)\end{aligned}$$

which can also be expressed as

$$\Big(\widetilde{F}'(x,y),\,\widetilde{G}'(x,y)\Big)=\int_0^\infty\frac{dk^2\,x^2(k^2)^{(1,2)}}{(k^2+1)(k^2+x)^4(k^2+y)}, \qquad (38)$$

which are always positive. There is no chargino contribution to C_2 and $\tilde{C}_2$ because of the color structure of the chargino box diagrams, and other terms are suppressed by the smallness of $\hat{Y}_q$. Since $\hat{Y}_t$ is large and $\hat{Y}_b$ can also be large for the case of large $\tan\beta$, we keep them in $C_{1,3}^{\tilde{\chi}^-}$. Note that in the AHS models, $\sum_{l,m}V^*_{lq}V_{mb}\delta^{lm}_{uLL}$ and $V^*_{tq}V_{tb}\delta^{33}_{uRR}$ are roughly of the order $|V_{tq}|$. As we will show soon, the Δm_{B_d} constraint require $\widetilde{m}\sim$ TeV due to large $\tilde{b}_R$–$\tilde{d}_R$ mixings. A typical LR stop mixing term contains $\hat{Y}_tV^*_{tq}\delta^{33}_{uLR}\sim 1.7\,V^*_{tq}m_t/\widetilde{m}$, which will be as small as $\sim 0.1V^*_{tq}$ for $\widetilde{m}\sim$ TeV. Furthermore, the flavor scale may not be too far from TeV [36], and there may not be much room for RG running to bring down the stop mass. Unlike the usual approach where one has light stop, the contributions from stop LR mixings are relatively small here.

<u>Gluino box</u> [41]:

$$\begin{aligned}C_1^{\tilde{g}}&=\frac{\alpha_s^2}{\widetilde{m}^2}\bigg[\frac{1}{4}\bigg(1-\frac{1}{N_c}\bigg)^2x_{\tilde{g}\tilde{q}}f_6(x_{\tilde{g}\tilde{q}})\\ &\qquad+\frac{1}{8}\bigg(N_c-\frac{2}{N_c}+\frac{1}{N_c^2}\bigg)\tilde{f}_6(x_{\tilde{g}\tilde{q}})\bigg](\delta^{q3}_{dLL})^2,\\ C_4^{\tilde{g}}&=\frac{\alpha_s^2}{\widetilde{m}^2}\bigg[\bigg(N_c-\frac{2}{N_c}\bigg)x_{\tilde{g}\tilde{q}}f_6(x_{\tilde{g}\tilde{q}})-\frac{\tilde{f}_6(x_{\tilde{g}\tilde{q}})}{N_c}\bigg]\delta^{q3}_{dLL}\delta^{q3}_{dRR},\\ C_5^{\tilde{g}}&=\frac{\alpha_s^2}{\widetilde{m}^2}\bigg[\frac{x_{\tilde{g}\tilde{q}}f_6(x_{\tilde{g}\tilde{q}})}{N_c^2}+\bigg(\frac{1}{2}+\frac{1}{2N_c^2}\bigg)\tilde{f}_6(x_{\tilde{g}\tilde{q}})\bigg]\delta^{q3}_{dLL}\delta^{q3}_{dRR},\end{aligned} \qquad (39)$$

where N_c is the number of colors, and

$$\begin{aligned}f_6(x)&=\frac{(17-9x-9x^2+x^3+6\ln x+18x\ln x)}{6(x-1)^5},\\ \tilde{f}_6(x)&=\frac{(1+9x-9x^2-x^3+6x\ln x+6x^2\ln x)}{3(x-1)^5}.\end{aligned} \qquad (40)$$

They are related to $\widetilde{F}'(x,y)$ and $\widetilde{G}'(x,y)$ by

$$\Big(xf_6(x),\,-\tilde{f}_6(x)\Big)=x^{-1}\Big(\widetilde{F}'(x^{-1},1),\,\widetilde{G}'(x^{-1},1)\Big). \qquad (41)$$

$\tilde{C}_1^{\tilde{g}}$ is obtained by interchanging $L\leftrightarrow R$ in $C_1^{\tilde{g}}$. We neglect $C^{\tilde{g}}_{2,3}$ and $\tilde{C}^{\tilde{g}}_{2,3}$ due to the smallness of LR and RL mixings. There are usual box and crossed box diagrams. Terms with N_c are from the former while $\mathcal{O}(1)$ terms are from the latter, which can be easily checked by 't Hooft double line notation. Terms with $1/N_c$, $1/N_c^2$ are subleading contributions from these two types of diagrams. We note that $C_4^{\tilde{g}}$ contains the largest N_c factor hence is the most sensitive to squark mixings. Note also that one has a zero in $C_1^{\tilde{g}}$ ($\tilde{C}_1^{\tilde{g}}$) for $x_{\tilde{g}\tilde{q}}\sim 2.43$.

<u>Gluino-neutralino box</u>:

$$\begin{aligned}C_1^{\tilde{g}\tilde{\chi}^0}&=\frac{\alpha_s\alpha_{\rm w}}{2\tilde{m}_{\tilde{g}}^2}\bigg(1-\frac{1}{N_c}\bigg)\bigg[G_L^jG_L^{*j}\tilde{G}'(x_{\tilde{q}\tilde{g}},x_{\tilde{\chi}^0_j\tilde{g}})\\ &\qquad-(G_L^jG_L^j+G_L^{*j}G_L^{*j})\sqrt{x_{\tilde{\chi}^0_j\tilde{g}}}\tilde{F}'(x_{\tilde{q}\tilde{g}},x_{\tilde{\chi}^0_j\tilde{g}})\bigg](\delta^{q3}_{dLL})^2,\\ C_2^{\tilde{g}\tilde{\chi}^0}&=\frac{\alpha_s\alpha_{\rm w}}{2m_{\tilde{g}}^2}\bigg(1-\frac{1}{N_c}\bigg)H^j_{bL}H^j_{bL}\\ &\qquad\times\sqrt{x_{\tilde{\chi}^0_j,\tilde{g}}}\tilde{F}'(x_{\tilde{q}\tilde{g}},x_{\tilde{\chi}^0_j\tilde{g}})\,(\delta^{q3}_{dLL})^2,\\ C_3^{\tilde{g}\tilde{\chi}^0}&=\frac{\alpha_s\alpha_{\rm w}}{2m_{\tilde{g}}^2}\bigg(1-\frac{1}{N_c}\bigg)H^j_{bL}H^j_{bL}\\ &\qquad\times\sqrt{x_{\tilde{\chi}^0_j,\tilde{g}}}\tilde{F}'(x_{\tilde{q}\tilde{g}},x_{\tilde{\chi}^0_j\tilde{g}})\,(\delta^{q3}_{dLL})^2,\\ C_4^{\tilde{g}\tilde{\chi}^0}&=\frac{\alpha_s\alpha_{\rm w}}{m_{\tilde{g}}^2}\bigg[\bigg(G_R^jG_L^j+G_R^{*j}G_L^{*j}-\frac{1}{N_c}H^j_{bL}H^j_{bR}\bigg)\\ &\qquad\times\tilde{G}'(x_{\tilde{q}\tilde{g}},x_{\tilde{\chi}^0_j\tilde{g}})-2\,(G_L^jG_R^{*j}+G_L^{*j}G_R^j)\\ &\qquad\times\sqrt{x_{\tilde{\chi}^0_j\tilde{g}}}\tilde{F}'(x_{\tilde{q}\tilde{g}},x_{\tilde{\chi}^0_j\tilde{g}})\bigg)\bigg]\delta^{q3}_{dLL}\delta^{q3}_{dRR},\\ C_5^{\tilde{g}\tilde{\chi}^0}&=\frac{\alpha_s\alpha_{\rm w}}{2\tilde{m}_{\tilde{g}}^2}\bigg[\bigg(H^j_{bL}H^j_{bR}-\frac{2}{N_c}G_R^jG_L^j-\frac{2}{N_c}G_R^{*j}G_L^{*j}\bigg)\\ &\qquad\times\tilde{G}'(x_{\tilde{q}\tilde{g}},x_{\tilde{\chi}^0_j\tilde{g}})+\frac{4}{N_c}(G_L^jG_R^{*j}+G_L^{*j}G_R^j)\\ &\qquad\times\sqrt{x_{\tilde{\chi}^0_j\tilde{g}}}\tilde{F}'(x_{\tilde{q}\tilde{g}},x_{\tilde{\chi}^0_j\tilde{g}})\bigg]\delta^{q3}_{dLL}\delta^{q3}_{dRR},\end{aligned} \qquad (42)$$

where j is summed over 1 to 4. The mixing matrices $G_{L,R},H_{bL,R}$ are given by

$$\begin{aligned}G_L^j&\equiv\tan\theta_WY_Q\mathcal{N}^*_{j1}+T_{3D}\mathcal{N}^*_{j2},\quad &G_R^j&\equiv\tan\theta_WQ_d\mathcal{N}^*_{j1}\\ H^j_{d,s,bL}&\equiv\mathcal{N}_{j3}\hat{Y}_{d,s,b}, &H^j_{d,s,bR}&\equiv\mathcal{N}^*_{j3}\hat{Y}_{d,s,b},\end{aligned} \qquad (43)$$

where Y_Q is the usual hypercharge, and the neutralino mixing matrix $\mathcal{N}$ diagonalizes the mass matrix $M_{\widetilde{\chi}^0}=\mathcal{N}^*\mathcal{M}\mathcal{N}^\dagger$,

$$\mathcal{M}=\begin{pmatrix}M_1 & 0 & -m_Zs_Wc_\beta & m_Zs_Ws_\beta\\ 0 & M_2 & m_Zc_Wc_\beta & -m_Zc_Ws_\beta\\ -m_Zs_Wc_\beta & m_Zc_Wc_\beta & 0 & -\mu\\ m_Zs_Ws_\beta & -m_Zc_Ws_\beta & -\mu & 0\end{pmatrix}.$$

$\tilde{C}^{\tilde{g}\tilde{\chi}^0}_{1,2,3}$ are obtained by chiral conjugation. Due to the smallness of $m_{d,s}$, we neglect terms with $H^j_{d,sL}$. Note that in Eq. (42), terms with G^jG^{*j} are from the usual box diagram, while terms with G^jG^j, $G^{*j}G^{*j}$, H^jH^j are from the crossed diagram. If the mass parameters M_1, M_2, $|\mu| \gg m_Z$, we will have simpler forms for $\mathcal{N}$, where $\mathcal{N}_{j1,j2} \sim \delta_{j1,j2}$, and higgsinos become maximally mixed. This will lead to a large cancellation of higgsino contributions in $C^{\tilde{g}\tilde{\chi}^0}_{2,3}$, $\tilde{C}^{\tilde{g}\tilde{\chi}^0}_{2,3}$. The diagonalization of the neutralino mass matrix usually leads to a negative mass eigenvalue, say $m_{\tilde{\chi}^0_i}$. It is well know that one can deal with it by two equivalent ways. One could either choose the phases in $\mathcal{N}$ such that $m_{\tilde{\chi}^0_i}$ are real and positive, or one could absorb the negative sign into $P_L\tilde{\chi}^0_i$, and modify Feynman rules accordingly [39,51]. However, there is a subtlety when dealing with the crossed diagrams in the latter approach. An additional negative sign is required for crossed box amplitudes when $\tilde{\chi}^0_i$ is in the loop, since $(\tilde{\chi}^0_i)^c = -\tilde{\chi}^0_i$ for that particular i.

Neutralino box :

$$
\begin{aligned}
C_1^{\tilde{\chi}^0} = & \frac{\alpha_w^2}{2m^2_{\tilde{\chi}^0_k}}\Big[G^j_LG^{*j}_LG^k_LG^{*k}_L\tilde{G}'(x_{\tilde{q}\tilde{\chi}^0_k},x_{\tilde{\chi}^0_j\tilde{\chi}^0_k}) \\
& -2\,G^j_LG^j_LG^{*k}_LG^{*k}_L\sqrt{x_{\tilde{\chi}^0_j\tilde{\chi}^0_k}}\tilde{F}'(x_{\tilde{q}\tilde{\chi}^0_k},x_{\tilde{\chi}^0_j\tilde{\chi}^0_k})\Big](\delta^{q3}_{dLL})^2, \\
C_2^{\tilde{\chi}^0} = & \frac{\alpha_w^2}{m^2_{\tilde{\chi}^0_k}}H^j_{bL}H^j_{bL}G^{*k}_LG^{*k}_L \\
& \times\sqrt{x_{\tilde{\chi}^0_j\tilde{\chi}^0_k}}\tilde{F}'(x_{\tilde{q}\tilde{\chi}^0_k},x_{\tilde{\chi}^0_j\tilde{\chi}^0_k})(\delta^{q3}_{dLL})^2, \\
C_3^{\tilde{\chi}^0} = & \frac{\alpha_w^2}{m^2_{\tilde{\chi}^0_k}}\sqrt{x_{\tilde{\chi}^0_j\tilde{\chi}^0_k}}\tilde{F}'(x_{\tilde{q}\tilde{\chi}^0_k},x_{\tilde{\chi}^0_j\tilde{\chi}^0_k})\big(H^j_{bL}H^j_{bL}G^{*k}_LG^{*k}_L \\
& -H^j_{bL}G^{*j}_LH^k_{bL}G^{*k}_L\big)(\delta^{q3}_{dLL})^2, \qquad (44)\\
C_4^{\tilde{\chi}^0} = & \frac{\alpha_w^2}{m^2_{\tilde{\chi}^0_k}}\tilde{G}'(x_{\tilde{q}\tilde{\chi}^0_k},x_{\tilde{\chi}^0_j\tilde{\chi}^0_k})\big(H^j_{bR}G^{*j}_LH^k_{bL}G^{*k}_R \\
& +H^j_{bL}H^j_{bR}G^{*k}_LG^{*k}_R\big)\delta^{q3}_{dLL}\delta^{q3}_{dRR}, \\
C_5^{\tilde{\chi}^0} = & \frac{2\alpha_w^2}{m^2_{\tilde{\chi}^0_k}}\Big[G^j_RG^j_LG^{*k}_RG^{*k}_L\tilde{G}'(x_{\tilde{q}\tilde{\chi}^0_k},x_{\tilde{\chi}^0_j\tilde{\chi}^0_k}) \\
& -2\,G^j_RG^{*j}_LG^k_LG^{*k}_R \\
& \times\sqrt{x_{\tilde{\chi}^0_j,\tilde{\chi}^0_k}}\tilde{F}'(x_{\tilde{q}\tilde{\chi}^0_k},x_{\tilde{\chi}^0_j\tilde{\chi}^0_k})\Big]\delta^{q3}_{dLL}\delta^{q3}_{dRR},
\end{aligned}
$$

where the indices j,k are summed over 1 to 4. We make use of the fact that $\tilde{G}'(x_{\tilde{q}\tilde{\chi}^0_k},x_{\tilde{\chi}^0_j\tilde{\chi}^0_k})/m^2_{\tilde{\chi}^0_k}$ and $\sqrt{x_{\tilde{\chi}^0_j,\tilde{\chi}^0_k}}\tilde{F}'(x_{\tilde{q}\tilde{\chi}^0_k},x_{\tilde{\chi}^0_j\tilde{\chi}^0_k})/m^2_{\tilde{\chi}^0_k}$ are symmetric under $j\leftrightarrow k$, which can be verified by using Eq. (38). $\tilde{C}^{\tilde{\chi}^0}_i$ are obtained by chiral conjugation. One can recognize contributions from the usual box or crossed box diagrams by similar rules stated earlier.

We note that the C_1s obtained in these four type of SUSY contributions are consistent with those in Ref. [49] by leading order Taylor expansion with respect to squark mixing angles.

After obtaining these Wilson coefficients at SUSY scale $M_{\rm SUSY}$, we apply renormalization group running to obtain B^0 mass scale values. The renormalization group running of these Wilson coefficients including leading order QCD corrections is given by [52],

$$
\begin{aligned}
C_1(\mu) &= \eta_1C_1(M_{\rm SUSY}), \\
C_2(\mu) &= \eta_{22}C_2(M_{\rm SUSY})+\eta_{23}C_3(M_{\rm SUSY}), \\
C_3(\mu) &= \eta_{32}C_2(M_{\rm SUSY})+\eta_{33}C_3(M_{\rm SUSY}), \qquad (45)\\
C_4(\mu) &= \eta_4C_4(M_{\rm SUSY})+\frac{1}{3}(\eta_4-\eta_5)C_5(M_{\rm SUSY}), \\
C_5(\mu) &= \eta_5C_5(M_{\rm SUSY}),
\end{aligned}
$$

where

$$
\begin{aligned}
&\eta_1 = \left(\frac{\alpha_s(M_{\rm SUSY})}{\alpha_s(m_t)}\right)^{6/21}\left(\frac{\alpha_s(m_t)}{\alpha_s(m_b)}\right)^{6/23}, \\
&\eta_2=\eta_1^{-2.42},\ \eta_3=\eta_1^{2.75},\ \eta_4=\eta_1^{-4},\ \eta_5=\eta_1^{1/2}, \\
&\eta_{22}=0.983\eta_2+0.017\eta_3,\ \eta_{23}=-0.258\eta_2+0.258\eta_3, \\
&\eta_{32}=-0.064\eta_2+0.064\eta_3,\ \eta_{33}=0.017\eta_2+0.983\eta_3. \quad (46)
\end{aligned}
$$

It is clear that $\eta_4 > \eta_2 > 1 > \eta_5 > \eta_1 > \eta_3$. Therefore, C_2 and, especially, C_4 will be enhanced by the RG running [52].

B. Impact on B_d Mixing

To obtain Δm_{B_d}, we use $\Delta m_{B_d} = 2|M^B_{12}|$, where

$$
\begin{aligned}
M^B_{12} &\equiv |M^B_{12}|\,e^{2i\Phi_{B_d}} \\
&= |M^{\rm SM}_{12}|\,e^{2i\phi_1}+|M^{\rm SUSY}_{12}|\,e^{i\phi_{\rm SUSY}}, \qquad (47)
\end{aligned}
$$

and

$$
M^{\rm SM}_{12} = 0.33\left(\frac{f_{B_d}\sqrt{\hat{B}_{B_d}}}{230\,{\rm MeV}}\right)^2\left(\frac{|V_{td}|}{8.8\times10^{-3}}\right)^2 e^{2i\phi_1}\ {\rm ps}^{-1}, \qquad (48)
$$

where $M^{\rm SM}_{12}$ is the SM contribution, its value is well known [53]. The vacuum insertion matrix elements of $\mathcal{O}_i$ are given in Ref. [41]. These matrix elements are modified by bag-factors to include non-factorizable effects. For simplicity, we assume the bag-factors for matrix elements of $\mathcal{O}_{2-5}$ are equal to $\hat{B}_{B_d}$, which is calculated for $\mathcal{O}_1$. In the subsequent numerical analysis, we take $f_{B_d}\hat{B}^{1/2}_{B_d} = (230\pm40)$ MeV [54]. For CKM matrix elements, we take $|V_{ub}/\lambda V_{cb}| = 0.41$ and $\phi_3 = 65°, 85°$. We use $|V_{td}|\times10^3 = 8.0,\ 9.2$ to get $\Delta m^{\rm SM}_{B_d}\sim 0.54,\ 0.72\,{\rm ps}^{-1}$, respectively, which are close to the experimental value $\Delta m_{B_d} = 0.484\pm0.010\,{\rm ps}^{-1}$ [12]. The large uncertainty in $f_{B_d}\hat{B}^{1/2}_{B_d}$ makes it possible for $\Delta m^{\rm SM}_{B_d}$ to lie within experimental range even for large ϕ_3 . However, when one considers $\Delta m_{B_s}/\Delta m_{B_d}$, the hadronic uncertainty is reduced

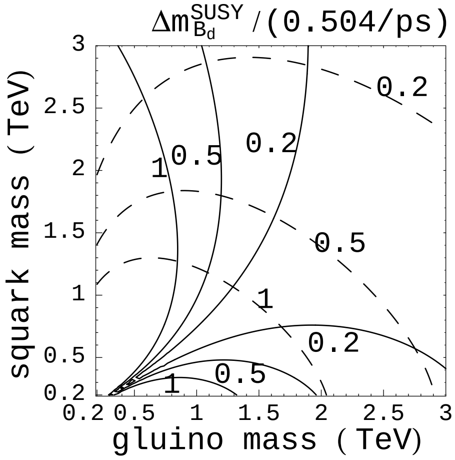

FIG. 2. Contribution to $(\Delta m_{B_d}^{\rm SUSY})$ / $(0.504\,{\rm ps}^{-1})$ from gluino box diagrams. The solid lines correspond to $\delta^{13}_{dRR} \sim 0.18$, while the dashed lines correspond to $\delta^{13}_{dLL}\delta^{13}_{dRR} \sim (0.18)^4$.

significantly, i.e. $\xi_s \equiv f_{B_s}\hat{B}_{B_s}^{1/2}/f_{B_d}\hat{B}_{B_d}^{1/2} = 1.16 \pm 0.05$ from lattice [55], and the SM prediction for large ϕ_3 case is not consistent with experiments, and New Physics would be needed for this case.

With the formulas above, we are ready to discuss the SUSY contributions to B–$\bar{B}$ mixing in the AHS models. In Fig. 2 we illustrate the $\widetilde{m}$–$m_{\tilde{g}}$ dependence of $\Delta m_{B_d}^{\rm SUSY}/(0.504\,{\rm ps}^{-1})$ from gluino box diagrams, where the denominator is the experimental bound at 2σ. The solid lines correspond to contributions from RR-RR mixings where each squark propagator has one $\delta^{13}_{dRR} \sim 0.18$ insertion. The dashed lines correspond to contributions from LL-RR mixings where one squark propagator has an insertion with $\delta^{13}_{dLL} \sim 0.18^3$ and the other an insertion with $\delta^{13}_{dRR} \sim 0.18$.

We can tell from Fig. 2 which mixings, LL-RR or RR-RR, give the dominant contribution in different regions of parameter space. For small $m_{\tilde{g}}$ the RR-RR mixings give tighter constraint, but for larger $m_{\tilde{g}}$ the LL-RR mixings is more stringent. The parameter space corresponding to $\Delta m_{B_d}^{\rm SUSY}/(0.504\,{\rm ps}^{-1}) \gg 1$ is excluded. Since contributions from other sparticles are sub-dominant in most of the parameter space, as we will discuss later, they are unlikely to cancel the gluino contribution.

We clearly need TeV range gluino and/or squark masses to satisfy the Δm_{B_d} constraint. This comes as a result of the large mixings in right-handed sector, and can be shown by simple arguments. $C_1^{\rm W}$ in the SM is roughly proportional to $(\alpha_{\rm W}/m_{\rm W}^2)^2 m_t^2 (V_{tb}V_{td}^*)^2$. For SUSY contribution assuming $m_{\tilde{g}} \sim \widetilde{m}$, the $(\alpha_{\rm W}/m_{\rm W}^2)^2 m_t^2$ factor is replaced by $N_c\alpha_S^2/\widetilde{m}^2$, and $(V_{tb}V_{td}^*)^2 \sim \lambda^6$ is replaced by $(\delta^{13}_{dRR})^2 \sim \lambda^2$ and $\delta^{13}_{dLL}\delta^{13}_{dRR} \sim \lambda^4$ for RR-RR and LL-RR mixing contributions, respectively. Since the SM contribution is already close to the experimental observation, one requires

$$\widetilde{m} \gtrsim \left(\frac{\alpha_S}{\alpha_{\rm W}}\right)\sqrt{\frac{N_c\lambda^{4,2}}{\lambda^6}}\left(\frac{m_{\rm W}^2}{m_t}\right) \sim 1,\, 4\,{\rm TeV}, \qquad (49)$$

from LL-RR and RR-RR mixings, respectively. Thus the typical scale of superparticles, $M_{\rm SUSY}$, has to be large due to large squark mixings. Comparing with Fig. 2, we note that the LL-RR case is close to this estimate, while $\tilde{m}$ for the RR-RR case is weaker than the estimate. This can be traced to the aforementioned cancellation in the RR-RR case where total cancellation in $\tilde{C}_1^{\tilde{g}}$ is possible for $x_{\tilde{g}\tilde{q}} \sim 2.43$. Considering RR-RR mixings only, there is a valley in the parameter space that even light superparticles with masses less than 250 GeV are allowed.

Neutralino box diagrams are induced by the same flavor source as the gluino box diagrams. The neutralino masses could be related to $m_{\tilde{g}}$ through a GUT-like relation on the gaugino Majorana masses [56],

$$m_{\tilde{g}} = \frac{\alpha_s}{\alpha_W}M_2, \qquad M_1 = \frac{5}{3}\frac{\alpha'}{\alpha_W}M_2. \qquad (50)$$

The gluino-neutralino box dominates over neutralino-neutralino box. The neutralino contribution to $\Delta m_{B_d}^{\rm SUSY}$ / $(0.504\,{\rm ps}^{-1})$ is less than 10% of gluino contribution for $m_{\tilde{g}}, \widetilde{m} > 500$ GeV for either RR-RR or LL-RR mixings with $\tan\beta = 2$–50 and $|\mu| = 100$–1000 GeV. For low $\widetilde{m}$, $|\mu|$ and large $\tan\beta$, $m_{\tilde{g}}$, its contribution, through $\tilde{C}_2^{\tilde{\chi}^0}$, can be comparable with the RR-RR mixing induced gluino contribution, which is suppressed by large $m_{\tilde{g}}$. However, as can be seen from Fig. 2, both the RR-RR and LL-RR mixing induced gluino contributions already give $\Delta m_{B_d}^{\rm SUSY}$ / $(0.504\,{\rm ps}^{-1}) \gg 1$ in that region, and we need to fine-tune the relative phase and size of these mixings to be within experimental bound.

The charged Higgs contributes $r \equiv |M_{12}^{\rm SUSY}/M_{12}^{\rm SM}| \sim$ (37, 11, 3)% for $m_{H^+} =$ (100, 400, 1000) GeV with $\tan\beta = 2$, and $\sim$ (0.3, 0.2, 0.1)% for $\tan\beta = 50$. The charged Higgs contribution interferes coherently with the SM amplitude. The result is consistent with early studies [50]. Without further interference from other SUSY contributions the $b \to s\gamma$ branching ratio constrains charged Higgs mass to be quite large [57]. In this case, the charged Higgs contribution to B^0–$\bar{B}^0$ mixing is small. It is known that cancellations from other sparticle contributions may reduce the bound on the charged Higgs mass [49,58]. We will return to this point in Section VI. On the other hand the chargino gives even smaller contribution, $r \sim 0.1$, 0.5% for TeV range $m_{\tilde{g}}$ and $\mu = \pm 1000$, ± 100 GeV, respectively. The main contribution comes from the term with $\hat{Y}_t^2\delta^{33}_{uRR}$ in $C_1^{\tilde{\chi}^-}$. The ratio is reduced by $\sim$ 30% for large $\tan\beta$ mainly due to the reduction of $\hat{Y}_t^2$. The smallness of chargino contributions is in contrast with early studies [50]. In our case there is no large mixing involving third generation in the up type sector, and there is no large splitting due to light stop. In this type of models, the SUSY contribution to B^0–$\bar{B}^0$ mixing is dominated by gluino exchange.

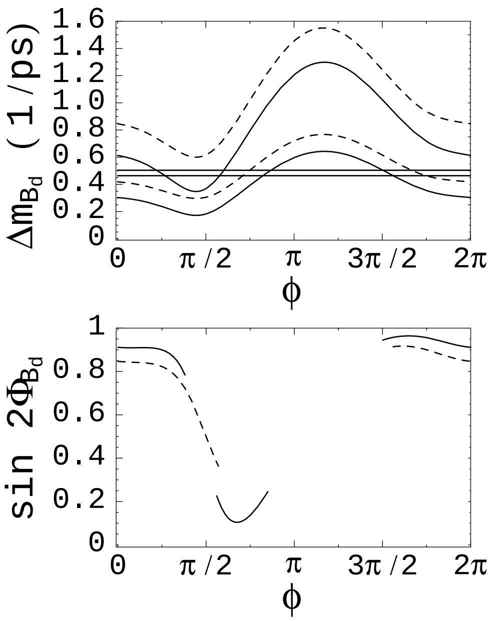

FIG. 3. One sigma range for (a) Δm_{B_d} and (b) $\sin 2\Phi_{B_d}$ vs. $\phi \equiv \arg \delta^{13}_{dRR}$, including both SM and SUSY effects, for gluino mass $m_{\tilde{g}} = 1.5$ TeV. The solid, short-dashed curves correspond to $\phi_3 = 65°$, $85°$ and $\widetilde{m} = 1.5$ TeV and $\tan\beta$ is taken as equal to 2. The horizontal lines in (a) indicate the 2σ experimental range.

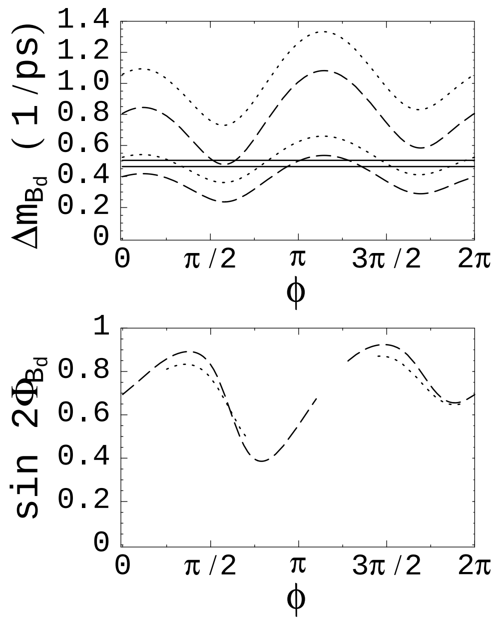

FIG. 4. Same as in Fig. 2 but for $\widetilde{m} = 3$ TeV, where long-dashed, dotted curves correspond to $\phi_3 = 65°$, $85°$.

After showing that gluino exchange gives dominant contributions to $M_{12}^{\rm SUSY}$, we turn to explore the interference between $M_{12}^{\rm SUSY}$ and $M_{12}^{\rm SM}$. We consider the SUSY phase $\phi \equiv \arg \delta^{13}_{dRR}$. The experimental measurement of Φ_{B_d} is no longer just ϕ_1 of SM. For illustration we plot, in Figs. 3(a) and (b), $\Delta m_{B_d} (\equiv 2|M^B_{12}|)$ and $\sin 2\Phi_{B_d}$ vs. ϕ, respectively, for 1.5 TeV squark mass and 1.5 TeV gluino mass. In Fig. 3(a), the solid (short-dashed) lines correspond to $\phi_3 = 65°\,(85°)$. The upper and lower solid and short-dashed lines denote the 1σ boundaries of $f_{B_d}\hat{B}^{1/2}_{B_d} = (230 \pm 40)$ MeV. If RR-RR mixings dominate, the SUSY phase $\phi_{\rm SUSY} \sim 2\phi$, while if LL-RR mixings dominate, then $\phi_{\rm SUSY} \sim \phi$. For $\phi_3 = 65°\,(85°)$, the SUSY model gives r ~22% (16%) from RR-RR mixings, and r ~55% (41%) from LL-RR mixings. These are consistent with Fig. 2, since $r \sim \Delta m^{\rm SUSY}_{B_d}/(0.504\,{\rm ps}^{-1})$. The two SUSY contributions interfere constructively (destructively) for $\phi \sim \pi$ (0). In the present case, LL-RR mixings dominate over RR-RR mixings and show a $\phi_{\rm SUSY} \sim \phi$ behavior in the graph.

In Fig. 4(a) and (b), we show the same physics measurable but with $\widetilde{m} = 3$ TeV, $m_{\tilde{g}} = 1.5$ TeV. For $\phi_3 = 65°\,(85°)$, the SUSY model contributes r ~29% (22%) from RR-RR mixings and r ~18%(13%) from LL-RR mixings vs $\Delta m^{\rm SM}_{B_d}$. These are again consistent with Fig. 2. The interference pattern is the same as the previous case. But in the present case, the RR-RR mixing contribution dominates over LL-RR mixings hence show a $\phi_{\rm SUSY} \sim 2\phi$ behavior in Fig. 4(a). We can see from Fig. 2 that RR-RR mixings dominate for the point $(m_{\tilde{g}}, \widetilde{m}) = (1.5,\, 3)$ TeV.

We see that $\sin 2\Phi_{B_d}$ as measured from $B_d \to J/\psi K_S$ can range from 0.1–0.95 and 0.4–0.9 as shown in Figs. 3(b) and 4(b), respectively. These curves are obtained by overlapping various segments from the corresponding curves within 1σ range of $f_{B_d}\hat{B}^{1/2}_{B_d}$. For example, in Fig. 3(a), any single line with $\widetilde{m} = m_{\tilde{g}} = 1.5$ TeV that corresponds to a value within 1σ of $f_{B_d}\hat{B}^{1/2}_{B_d}$ should lie within the two solid lines. Each single line only intercepts with the 2σ experimental range for Δm_{B_d} for some region of ϕ and corresponds to some segment in the solid line of Fig. 3(b) (one can compare with Fig. 1(b) of Ref. [46], where $f_{B_d}\hat{B}^{1/2}_{B_d} = 200$ MeV is used). Taking the uncertainty of $f_{B_d}\hat{B}^{1/2}_{B_d}$ into account enlarges the parameter space considerably. The fact that these segments from different value of $f_{B_d}\hat{B}^{1/2}_{B_d}$ lie on a line and not forming a band corresponds to our simplifying assumption of using same bag-factor for all $\mathcal{O}_i$. Thus, they can be factored out without affecting the argument of M_{12} in Eq. (47).

In SM, we have $\sin 2\phi_1 \simeq 0.75$–0.71 [10] for $\phi_3 = 65°$–$85°$. The measurements from Belle [6] and BaBar [7] in year 2000 indicated smaller values for $\sin 2\phi_1$ vs SM, and would correspond to relatively specific $\phi \equiv \arg \delta^{13}_{dRR}$ phase values (between $\pi/2$ and π) in Figs. 3(b) and 4(b). More recent definitive measurements from BaBar [8] and Belle [9], in year 2001, give $\sin 2\phi_1$ values close to 1! The average of BaBar and Belle 2001 values is given in Eq. (2). It is rather intriguing that this range corresponds to the main parameter space allowed by B_d mixing, suggesting that SUSY contributions could be comparable to SM. In particular, if the Belle 2001 central value of $\sin 2\phi_1 = 0.99$ holds up, it would imply that $\phi \equiv \arg \delta^{13}_{dRR}$ is between

$3\pi/2$ and 2π. The lighter squark mass case of Figs. 3 is preferred, but the heavier squark mass case of Figs. 4 is also possible.

To conclude this section, we note that one requires heavy squark and gluino masses to satisfy Δm_{B_d} bound. While the direct search of such heavy superparticles become less promising, these particles can show their effect in B_d–$\bar{B}_d$ mixing phases, even with TeV masses.

C. Brief Discussion on B_s Mixing and CP Phase

As shown in previous section, one could have s flavor decoupled and the above discussion is applicable, and is being tested right now. Alternatively, and mutually exclusive to the above case, it could be the d flavor that is decoupled, and the B_d system would be SM-like, which may still turn out to be the case in 2002. If so, B_s mixing may be the place where SUSY AHS effects show up.

With the CKM like relation $\delta^{13}_{dRR}/\delta^{23}_{dRR} \sim V_{td}/V_{ts} \sim \lambda$ from Eqs. (17), (18), the B_s mixing case is rather similar to the discussion of B_d mixing, so long that the mass insertion approximation can be used. One can just scale up from previous discussion. The gluino contribution to B^0_s–$\bar{B}^0_s$ mixing is discussed in Ref. [47] and [48], where in the latter work the mass insertion approximation is relaxed. Since $\tilde{s}_R$–$\tilde{b}_R$ mixing ~ 1 in SUSY AHS model, one in principle could have a relatively light "strange-beauty" squark, which, unlike the heavy SUSY scale that is the focus of this paper, can impact on direct search. We will discuss this case later in Sec. VII.

The B_s mixing phase may not be vanishingly small as in SM, and can be searched for at the Tevatron collider in a matter of years. The most interesting situation would be to find (soon!) Δm_{B_s} not far above SM expectation, but with large $\sin 2\Phi_{B_s}$.

IV. CONSTRAINTS FROM K^0–$\bar{K}^0$ MIXING

As shown in Eqs. (14) and (17), AHS models not only give large mixing in RR sector involving third generation down squark, they also give large mixing in 1-2 generations. It is well known that Δm_K is much smaller than Δm_{B_d} hence offers a much stronger constraint, while ε_K is even tighter. They make $\delta^{12}_{dLL,RR} \sim \lambda$ impossible to sustain even with $\widetilde{m}$, $m_{\tilde{g}} \gtrsim$ TeV [42,43]. The formulas for kaon mixing are similar to that for B_q–$\bar{B}_q$ mixing, with only some modifications needed. For charged Higgs exchange diagrams, Eq. (33) now becomes,

$$\begin{aligned} C_1^{H^-} &= \frac{\alpha_{\rm W}^2}{8m_W^2} V_{is}V_{js}V^*_{id}V^*_{jd}[x_{i{\rm W}}\, x_{jH} \cot^4\beta \frac{1}{4} G(x_{iH}, x_{jH}) \\ &\quad +2x_{i{\rm W}}x_{j{\rm W}}\cot^2\beta (F'(x_{i{\rm W}}, x_{j{\rm W}}, x_{H{\rm W}}) \\ &\quad +\frac{1}{4}G'(x_{i{\rm W}}, x_{j{\rm W}}, x_{H{\rm W}})], \end{aligned} \tag{51}$$

where i, j are generation indices of up type quarks and summed over. Other terms are neglected due to the smallness of quark masses $m_{s,d}$. Chargino contributions are modified by changing $V^*_{lq}V_{mb}\delta^{lm}_{uLL}$ in Eq. (34) to $V^*_{ld}V_{ms}\delta^{lm}_{uLL}$ and neglecting other terms. Gluino and neutralino contributions are modified by changing δ^{q3} in Eqs. (39), (42), (44) to δ^{12} and neglecting all $H_{L,R}$ terms. The QCD running formula is also modified accordingly.

The charged Higgs contributions are in general small. For gluino and neutralino contributions we show in Table II the limits on $\sqrt{|{\rm Re}\,\delta^{12}_{AB}\delta^{12}_{CD}|}$ and $\sqrt{|{\rm Im}\,\delta^{12}_{AB}\delta^{12}_{CD}|}$ from $(\Delta m_K^{\rm SUSY}) < 3.521\times 10^{-12}$ MeV for $\widetilde{m} = 1.5$ TeV, where $A, B, C, D = L, R$. The constraints from ε_K can also be estimated by using

$$|\varepsilon_K| = \frac{|{\rm Im}\, M_{12}|}{\sqrt{2}\Delta m_K} < 2.268\times 10^{-3}. \tag{52}$$

For different values of $\widetilde{m}$, the limits can be roughly obtained by multiplying a factor $\widetilde{m}/(1.5\,{\rm TeV})$. We take the hadronic scale in Eq. (46) to be $\sim$ GeV and $\alpha_s(M_{\rm Z}) = 0.1185$. We can reproduce the results of Ref. [52] by using $\alpha_s(\mu) \sim 1$ and $\widetilde{m} = 500$ GeV and by considering gluino contributions only. The QCD effects enhance the SUSY contributions by a few times for contributions arising from $\delta_{dLL}\,\delta_{dRR}$, as one can see from $\eta_4 \sim 6$. Since the most severe constraint is on $\delta_{dLL}\delta_{dRR}$, the QCD effects make it more stringent [52]. From Eq. (52), ε_K gives even more stringent constraint than Δm_K. However, the constraint becomes less severe if the phases of δs are small (of order 0.01). Together with constraint from electric dipole moment of neutron, one may be led to the idea of approximate CP (see for example [59]).

From the above discussion, it is clear that δ^{12}_{dLL}, $\delta^{12}_{dRR} \sim \lambda$ cannot be sustained. One has to invoke QSA as discussed in Sec. II to impose appropriate "texture zeros". We see that the values given in Table I are all well below the limits from Δm_K and ε constraints, Table II, even with $\mathcal{O}(1)$ phases. The approximate CP assumption can be relaxed.

x	$\sqrt{\|{\rm Re}\,(\delta^{12}_{dLL})^2\|}$	$\sqrt{\|{\rm Re}\,\delta^{12}_{dLL}\delta^{12}_{dRR}\|}$	$\sqrt{\|{\rm Re}\,(\delta^{12}_{dRR})^2\|}$
0.3	8.9×10^{-2}	4.3×10^{-3}	9.3×10^{-2}
1	2.3×10^{-1}	4.9×10^{-3}	2.0×10^{-1}
4	2.5×10^{-1}	7.0×10^{-3}	4.3×10^{-2}
	$\sqrt{\|{\rm Im}\,(\delta^{12}_{dLL})^2\|}$	$\sqrt{\|{\rm Im}\,\delta^{12}_{dLL}\delta^{12}_{dRR}\|}$	$\sqrt{\|{\rm Im}\,(\delta^{12}_{dRR})^2\|}$
0.3	7.1×10^{-3}	3.4×10^{-4}	7.4×10^{-3}
1	1.8×10^{-2}	3.9×10^{-4}	1.6×10^{-2}
4	2.0×10^{-2}	5.6×10^{-4}	3.4×10^{-2}

TABLE II. Limits on $\sqrt{|{\rm Re}\,\delta^{12}_d\,\delta^{12}_d|}$ and $\sqrt{|{\rm Im}\,\delta^{12}_d\,\delta^{12}_d|}$ for squark mass $\widetilde{m} = 1.5$ TeV and for different $x = m^2_{\tilde{g}}/\widetilde{m}^2$, including leading order QCD corrections. The constraints are from $\Delta m_K^{\rm SUSY} < 3.521\times 10^{-12}$ MeV and $|\varepsilon_K^{\rm SUSY}| < 2.268\times 10^{-3}$.

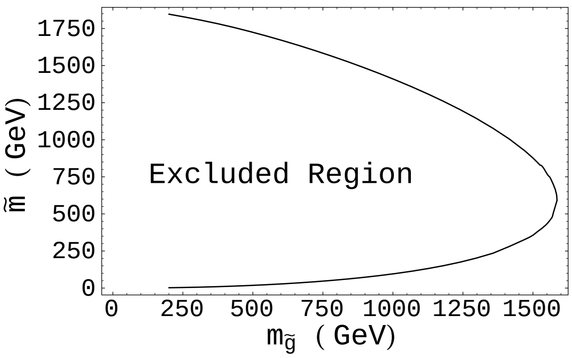

FIG. 5. Limit on $\widetilde{m}$ vs $m_{\tilde{g}}$ from chargino contributions to Δm_K arising from $\delta^{12}_{uLL} \sim \lambda$.

As noted in Sec. II, QSA will induce $\delta^{12}_{uLL} \sim \lambda$ by shifting the source of the Cabibbo angle to the up-type sector. The full strength δ^{12}_{uLL} can contribute to kaon mixing via chargino diagrams. In Fig. 5, we show the parameter space constrained by Δm_K. We use the GUT relation on gaugino mass as given in Eq. (50) for sake of simplicity and definiteness. The horizontal axis can be converted to wino mass by multiplying $m_{\tilde{g}}$ by ~ 0.4. For the ε_K constraint, we would need $\arg(\delta_u)$ to be less than 0.1. We see from Fig. 5 that the kaon mixing constraint also points to TeV scale gluino and squarks. We stress that this is a generic feature of QSA and has nothing to do with the choice of retaining M_d^{3i} or not. Keeping M_d^{3i} leads to interesting low energy physics, even with TeV scale particles as a result of kaon mixing constraints.

We note that it is possible for the chargino diagrams to interfere destructively with LL (or RR) mixing induced gluino and neutralino contributions, and one can satisfy the kaon constraint with a lower mass scale. However, since these correspond to different set of parameters, it is unlikely for the interference to be destructive in general.

V. IMPLICATIONS FOR D^0–$\bar{D}^0$ MIXING

The experimental situation for D^0–$\bar{D}^0$ mixing is rather volatile at the present time. A search by the CLEO Collaboration gives $1/2x'^2_D < 0.041\%$ and $-5.8\% < y'_D < 1.0\%$ [15], where x'_D and y'_D are defined in Eq. (5). CLEO further adopted $\delta_D \simeq 0$ from model arguments to reach a more stringent bound of $x_D \simeq x'_D < 2.9\%$. If $\delta_D \neq 0$ [25], however, the preferred negative value of y'_D may in fact be hinting at $x_D \sim$ the few % level.

Another approach is to compare $D^0 \to K^-\pi^+$ and K^-K^+ decays and measure the lifetime difference between CP even and odd final states. The current world average from Belle [18], BaBar [19], CLEO [17], E791 [14] and FOCUS [16] Collaborations is $(1.1 \pm 0.99)\%$ [22]. Since this is consistent with zero and does not support the nonzero claim by FOCUS, we shall take the more stringent constraint of $x_D < 2.9\%$ from CLEO in the following. What we find intriguing, however, is that $\delta^{12}_{uLL} \sim \lambda$

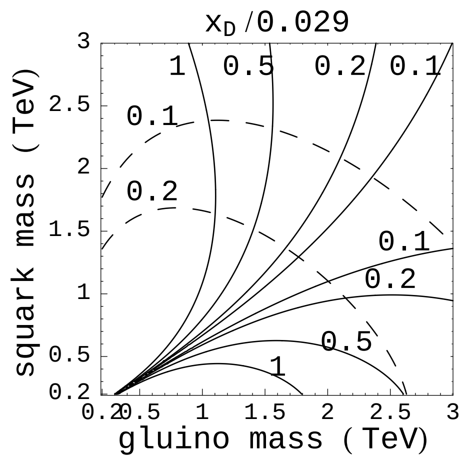

FIG. 6. Contribution to $x_D/2.9\%$ from gluino box diagrams. The solid lines correspond to $\delta^{12}_{uLL} \sim 0.18$, while the dashed lines correspond to $\delta^{12}_{uLL}\delta^{12}_{uRR} \sim (0.18)^{5.5}$.

with $\widetilde{m}$, $m_{\tilde{g}} \sim$ TeV brings x_D right into the ball-park of the % level! Furthermore, this can be probed in detail in the next few years at the B factories, and in the longer run, by hadron collider detectors.

We consider gluino and neutralino exchange diagrams induced by up-squark mixing with $\mu = 1$ TeV and $\tan\beta = 2$ with formulas similar to B mixing. The dependence on $\tan\beta$ is weak. The SUSY contribution from gluino box to $x_D/0.029$ is illustrated in Fig. 6 in the $m_{\tilde{g}}$–$\widetilde{m}$ plane. The solid lines correspond to $\delta^{12}_{uLL} \sim 0.18$, while the dashed lines correspond to $\delta^{12}_{uLL}\delta^{12}_{uRR} \sim (0.18)^{5.5}$. It is clear that LL-LL induced gluino box diagrams dominate. In Fig. 7 we illustrate x_D vs. $\widetilde{m}$ for $m_{\tilde{g}} = 0.8$, 1.5, 3 TeV, respectively. As in the B mixing case, there is a narrow valley from δ_{uLL} induced gluino contributions around $m^2_{\tilde{g}}/\widetilde{m}^2 \sim 2.43$ when $C^{\tilde{g}}_1$ of Eq. (39) vanishes. This could make the parameter space from Fig. 6 too restrictive. However, the actual zeros in Fig. 7, occur at slightly shifted mass ratios, reflecting a cancellation between various contributions from δ_{uLL} and $\delta_{uLL}\delta_{uRR}$ when they have a common phase. Although ε_K constrains $\arg(\delta^{12}_{uLL})$ to be less than 0.1, the phase of δ^{12}_{uRR}

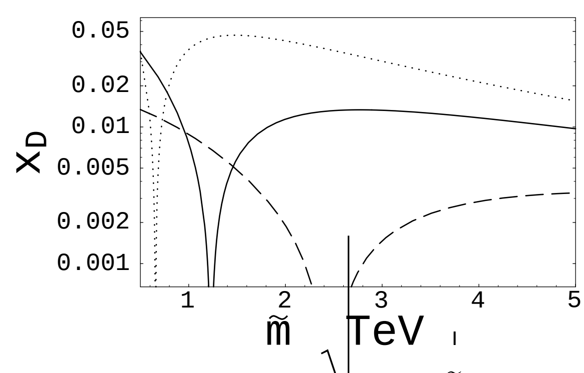

FIG. 7. Gluino contribution to x_D vs. $\widetilde{m}$. Dotted, solid and dashed lines are for $m_{\tilde{g}} = 0.8$, 1.5 and 3 TeV, respectively.

is not constrainted since δ_{uRR}^{12} is by itself small. In general, the SUSY phase δ_{uRR}^{12} does not have to vanish, and having phase in common with δ_{uLL} is not likely. Thus, the deep valley would in general be filled, but the figure illustrates the adjustability of x_D. It also gives an explicit example where detectable D^0 mixing would likely [23] carry a CP violating phase.

To conclude this section, we note that AHS with QSA is known to produce large D meson mixings. The stringent upper bound on x_D seem to provide severe constraint for QSA models [36,60]. This is more or less true when squarks and gluino are as light as a few hundred GeV. However, as a result of the large mixing in squark sector in our case, the proximity of Δm_{B_d} to SM expectation leads to squarks at TeV scale, and $\sin 2\phi_1$ may be affected in an interesting way. It is intersting that the scale determined from this leads to a D^0 meson mixing close to experimental hints. We eagerly await the experimental situation to clear up, i.e. whether the CLEO hint is due to $\Delta\Gamma_D$ [22] or Δm_D.

VI. RADIATIVE B DECAYS

The effective Hamiltonian for $b \to q\gamma$, qg transitions, where $q = d$ or s, is given by

$$H_{\rm eff.} = -\frac{G_F}{\sqrt{2}} \frac{m_b}{4\pi^2} V_{tb} V_{tq}^* \bar{q} \left\{ e \left[C_{7\gamma}\, R + C'_{7\gamma}\, L \right] F^{\mu\nu} + g \left[C_{8g}\, R + C'_{8g}\, L \right] T^a G_a^{\mu\nu} \right\} \sigma_{\mu\nu} b, \tag{53}$$

where we have neglected m_q, $C_{7,8} = C_{7,8}^{\rm SM} + C_{7,8}^{\rm New}$ are the sum of SM and New Physics contributions, while $C'_{7,8}$ come purely from New Physics. We are particularly interested in the case where $C'_{7\gamma,8g}$ are large. The effects from the SUSY contributions are given by

$$C_{7\gamma}^{(\prime){\rm New}} = C_{7\gamma,H^-}^{(\prime)} + C_{7\gamma,\tilde{g}}^{(\prime)} + C_{7\gamma,\tilde{\chi}^-}^{(\prime)} + C_{7\gamma,\tilde{\chi}^0}^{(\prime)}, \tag{54}$$

$$C_{8g}^{(\prime){\rm New}} = C_{8g,H^-}^{(\prime)} + C_{8g,\tilde{g}}^{(\prime)} + C_{8g,\tilde{\chi}^-}^{(\prime)} + C_{8g,\tilde{\chi}^0}^{(\prime)}. \tag{55}$$

The Feynman diagrams are shown in Fig. 8.

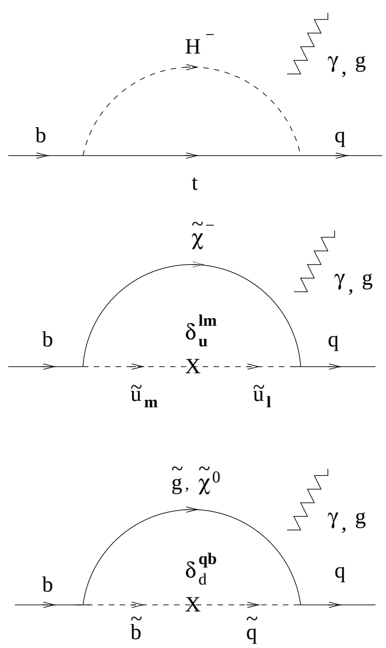

FIG. 8. SUSY penguin diagrams for $b \to q\gamma, qg$ processes.

A. Formulas

Charged Higgs Exchange:

$$C_{7\gamma,H^-} = -\frac{x_{tH}}{2} \Big\{ \cot^2\beta \left[Q_u F_1(x_{tH}) + F_2(x_{tH}) \right] + \left[Q_u F_3(x_{tH}) + F_4(x_{tH}) \right] \Big\}, \tag{56}$$

$$C_{8g,H^-} = -\frac{x_{tH}}{2} \Big[\cot^2\beta F_1(x_{tH}) + F_3(x_{tH}) \Big], \tag{57}$$

where $F_i(x)$ are loop functions and the explicit expressions can be found in Ref. [49].

Gluino Exchange:

$$C_{7\gamma,\tilde{g}} = \frac{\pi\alpha_s}{\sqrt{2} G_F V_{tb} V_{td}^*} \frac{Q_d 2\, C_2(R)}{\tilde{m}^2} \times \Big\{ \delta_{dLL}^{13}\, g_2(x_{\tilde{g}\tilde{q}}) - \frac{m_{\tilde{g}}}{m_b} \delta_{dLR}^{13}\, g_4(x_{\tilde{g}\tilde{q}}) \Big\}, \tag{58}$$

$$C_{8g,\tilde{g}} = \frac{\pi\alpha_s}{\sqrt{2} G_F \tilde{m}^2 V_{tb} V_{td}^*} \times \Big\{ \delta_{dLL}^{13} \Big[\left[2C_2(R) - C_2(G) \right] g_2(x_{\tilde{g}\tilde{q}}) - C_2(G)\, g_1(x_{\tilde{g}\tilde{q}}) \Big] + \frac{m_{\tilde{g}}}{m_b} \delta_{dLR}^{13} \Big[\left[C_2(G) - 2C_2(R) \right] g_4(x_{\tilde{g}\tilde{q}}) + C_2(G)\, g_3(x_{\tilde{g}\tilde{q}}) \Big] \Big\}, \tag{59}$$

where Q_d is the down quark electric charge, $C_2(G) = N = 3$ and $C_2(R) = (N^2-1)/(2N) = 4/3$ are Casimirs, and the functions $g_i(x) = -d/dx[x\, F_i(x)]$, i.e.

$$\begin{aligned} g_1(x) &= \frac{1 + 9x - 9x^2 - x^3 + 6x(1+x)\, \ln x}{6\,(x-1)^5}, \\ g_2(x) &= \frac{1 - 9x - 9x^2 + 17x^3 - 6x^2(3+x)\, \ln x}{12\,(x-1)^5}, \\ g_3(x) &= \frac{5 - 4x - x^2 + (2+4x)\, \ln x}{2\,(x-1)^4}, \\ g_4(x) &= \frac{-1 - 4x + 5x^2 - 2x(2+x)\, \ln x}{2\,(x-1)^4}. \end{aligned} \tag{60}$$

The chirality partners $C'_{7\gamma,8g}$ are obtained by interchanging L and R in the δ's. We see that δ_{LL} and δ_{LR} contribute to $C_{7\gamma,8g}$, while δ_{RR} and δ_{RL} contribute to $C'_{7\gamma,8g}$.

There is an enhancement factor $m_{\tilde g}/m_b$ that comes with $\delta_{LR,RL}$. The factor m_b is from normalizing with respect to the SM result, and hence the smallness of the b quark mass with respect to the gluino mass is the origin of this enhancement. Such enhancement was noted in our earlier study [30] of $\tilde s$-$\tilde b$ mixings where $\tilde d$ sector was decoupled completely. It has also been invoked to generate ε'/ε via an analogous δ^{12}_{LR} term [34,61] under a horizontal U(2) (hence non-Abelian) symmetry model. The mechanism is generic and has been discussed in Ref. [62], but SUSY with LR squark mixings gives a beautiful example.

<u>Chargino Exchange:</u>

$$
\begin{aligned}
&C_{7\gamma,\tilde\chi^-} = \frac{m_w^2}{\widetilde m^2 V_{tb}V_{tq}^*}\\
&\times\Bigg\{\Bigg[\mathcal V_{j1}\mathcal V_{j1}^* V_{lq}^* V_{mb}\delta^{lm}_{uLL} - \mathcal V_{j1}\mathcal V_{j2}^*\hat Y_t V_{lq}^* V_{tb}\delta^{lt}_{uLR}\\
&\quad -\mathcal V_{j2}\mathcal V_{j1}^*\hat Y_t V_{tq}^* V_{mb}\delta^{tm}_{uRL} + \mathcal V_{j2}\mathcal V_{j2}^*\hat Y_t^2 V_{tq}^* V_{tb}\delta^{tt}_{uRR}\Bigg]\\
&\times\Big[g_1(x_{\tilde\chi_j^-\tilde q}) + Q_u g_2(x_{\tilde\chi_j^-\tilde q})\Big] - \frac{m_{\tilde\chi_j^-}}{m_b}\Big[\mathcal V_{j1}\mathcal U_{j2}\hat Y_b V_{lq}^* V_{mb}\delta^{lm}_{uLL}\\
&\quad -\mathcal V_{j2}\mathcal U_{j2}\hat Y_t\hat Y_b V_{tq}^* V_{mb}\delta^{tm}_{uRL}\Big]\Big[g_3(x_{\tilde\chi_j^-\tilde q}) + Q_u g_4(x_{\tilde\chi_j^-\tilde q})\Big]\Bigg\},
\end{aligned}
\tag{61}
$$

$$
\begin{aligned}
&C'_{7\gamma,\tilde\chi^-} = \hat Y_q\,\frac{m_w^2}{\widetilde m^2 V_{tb}V_{tq}^*}\\
&\times\Big\{\mathcal U_{j2}\mathcal U_{j2}^*\hat Y_b V_{lq}^* V_{mb}\delta^{lm}_{uLL}\Big[g_1(x_{\tilde\chi_j^-\tilde q}) + Q_u g_2(x_{\tilde\chi_j^-\tilde q})\Big]\\
&\quad +\frac{m_{\tilde\chi_j^-}}{m_b}\Big[g_3(x_{\tilde\chi_j^-\tilde q}) + Q_u g_4(x_{\tilde\chi_j^-\tilde q})\Big]\\
&\times\Big[\mathcal V_{j2}^*\mathcal U_{j2}^*\hat Y_t V_{lq}^* V_{tb}\delta^{lt}_{uLR} - \mathcal U_{j2}^*\mathcal V_{j1}^* V_{lq}^* V_{mb}\delta^{lm}_{uLL}\Big]\Big\},
\end{aligned}
\tag{62}
$$

where as before we sum over l, m for three generations and j for two chargino mass eigenstates. $C^{(\prime)}_{8g,\tilde\chi^-}$ can be obtained by dropping $g_{1,3}$ from above equations and replacing Q_u with 1. It is clear from these equations that $C'_{7\gamma,\tilde\chi^-}$ is suppressed by $\hat Y_q$.

<u>Neutralino Exchange:</u>

$$
\begin{aligned}
C_{7\gamma,\tilde\chi^0} &= \frac{Q_d m_w^2}{\widetilde m^2 V_{tb}V_{tq}^*}\Big\{2G^{*j}_{qL}G^j_{bL}\delta^{i3}_{LL}g_2(x_{\tilde\chi_j^0\tilde q})\\
&\quad +\frac{m_{\tilde\chi_j^0}}{m_b}\Big[-2G^{*j}_{qL}G^j_{bR}\delta^{i3}_{dLR} + \sqrt2 G^{*j}_{qL}H^j_{bL}\delta^{i3}_{dLL}\Big]g_4(x_{\tilde\chi_j^0\tilde q})\Big\},
\end{aligned}
\tag{63}
$$

where j is summed over four neutralino mass eigenstates, and $C^{(\prime)}_{8g,\tilde\chi^0}$ can be obtained by replacing $Q_d \to 1$ in the above equation. We neglect terms with $H^*_{qL,R}$ and some LR mixing terms when there is no chiral enhancement.

Similar to the gluino case, the chiral partners C' are obtained by taking a conjugation in the chirality $L \leftrightarrow R$ and noting that $G_{L,R} \leftrightarrow -G_{R,L}$.

When running down to the B decay scale $\mu \approx m_b$, the leading order Wilson coefficients $C_i^{(\prime)}$ are given by [63],

$$
\begin{aligned}
C_{7\gamma}(\mu = m_b) &= -0.31 + \eta_7\, C^{\rm New}_{7\gamma}(M_{\rm SUSY})\\
&\qquad +\frac{8}{3}(\eta_8 - \eta_7)\, C^{\rm New}_{8g}(M_{\rm SUSY}),\\
C_{8g}(\mu = m_b) &= -0.15 + \eta_8\, C^{\rm New}_{8g}(M_{\rm SUSY}),
\end{aligned}
\tag{64}
$$

where,

$$
\begin{aligned}
\eta_7 &= \left(\frac{\alpha_s(M_{\rm SUSY})}{\alpha_s(m_t)}\right)^{16/21}\left(\frac{\alpha_s(m_t)}{\alpha_s(m_b)}\right)^{16/23},\\
\eta_8 &= \left(\frac{\alpha_s(M_{\rm SUSY})}{\alpha_s(m_t)}\right)^{14/21}\left(\frac{\alpha_s(m_t)}{\alpha_s(m_b)}\right)^{14/23},
\end{aligned}
\tag{65}
$$

while for opposite chirality, which receives no SM contribution, one simply replaces $C^{\rm New}$ by C' and set the constant terms to zero.

B. Phenomenological Impact

It is well known that $B \to X_s\gamma$ is a severe constraint on New Physics. The current experimental results are $Br(B \to X_s\gamma) = (2.85 \pm 0.35 \pm 0.22) \times 10^{-4}$ [64], $(3.37 \pm 0.53 \pm 0.42^{+0.50}_{-0.54}|_{\rm model}) \times 10^{-4}$ [65], $(3.11 \pm 0.82 \pm 0.72) \times 10^{-4}$ [66], from CLEO, Belle and ALEPH, respectively. It is known that the charged Higgs contribution interferes constructively [57] with SM contribution $C^{\rm SM}_{7\gamma}(m_b) = -0.31$. By using Eq.(56), we have,

$$
\frac{C_{7\gamma,H^-}}{C^{\rm SM}_{7\gamma}} = 1 + (35\%, 22\%, 15\%, 11\%, 9\%),
\tag{66}
$$

for $m_{H^-} = (400, 600, 800, 1000, 1200)$ GeV. The rate can get enhanced by $\sim$ 80%–20%. The result is insensitive to $\tan\beta$ within 2–50, since the term without $\cot\beta$ in Eq. (56) is dominant [57]. If we require that the deviation from the SM rate to be less than 20%, which is close to experimental error, we need $m_{H^+} \geq 1.2$ TeV, or we may need cancellations from other particles [49]. The chargino contribution may partially cancel the charged Higgs contribution [58]. This mechanism is still operative even if we have $m_{\tilde g}, \widetilde m = 1.5$ TeV.

$$
\begin{aligned}
\frac{C_{7\gamma,\tilde\chi^-}}{C^{\rm SM}_{7\gamma}} &= 1 \pm \left\{\begin{matrix}13\ (0.6)\\ 17\ (0.9)\\ 12\ (0.6)\end{matrix}\right\}\%\,\frac{\delta^{33}_{uRL}}{m_t/\widetilde m}\\
&\quad +\left\{\begin{matrix}18\ (0.5)\\ 8\ (0.05)\\ 3\ (0.08)\end{matrix}\right\}\%\,\frac{V_{ls}^* V_{mb}\delta^{lm}_{uLL}}{\lambda^2}\\
&\quad -\left\{\begin{matrix}0.3\ (0.4)\\ 0.2\ (0.2)\\ 0.1\ (0.1)\end{matrix}\right\}\%\,\frac{V_{ts}^* V_{tb}\delta^{33}_{uRR}}{\lambda^2},
\end{aligned}
\tag{67}
$$

for $\mu = \pm(100, 500, 1000)$ GeV, and $\tan\beta = 50$ (2). The δ^{33}_{uRL} term comes from chiral enhancement (with $m_{\tilde{\chi}^-_j}/m_b$ factor). The sign of the coefficient in front of δ^{33}_{uRL} is the same as the sign of μ. The coefficient for $|\mu| = 500$ GeV is greater than $|\mu| = 100$ GeV, because of chiral enhancement. The coefficient will drop below 0.1 for $|\mu| \geq 1.2$ TeV. LR and RL mixings without chiral enhancement are negligible, their contributions being only about $10^{-5,-6}$ of $C^{\rm SM}_{7\gamma}$. The term with $V^*_{ls} V_{mb} \delta^{lm}_{uLL}$ is also from the chiral enhancement term, while the last term is not chirally enhanced. Due to the smallness of $\hat{Y}_s$, $|C'_{7\gamma,\tilde{\chi}^-}/C^{\rm SM}_{7\gamma}|$ is below 1% for expected mixing angles . For $\tan\beta = 2$, terms are suppressed by the $\hat{Y}_b(\tan\beta = 2)/\hat{Y}_b(\tan\beta = 50) \sim 1/22$ factor, except for the last term.

The sign of RL stop mixings is anti-correlated to μ. This can provide needed cancellations for low m_{H^-}, even if μ, M_2 and $\widetilde{m}$ are large. For $|\mu| = 100, 500$ GeV, if the signs of second and third terms of Eq. (67) are negative, the charged Higgs mass can be as low as 300, 400 GeV, for a rate within 20% from SM expectation. Even for $|\mu| = 1$ TeV, the cancellation may lower m_{H^-} to 600 GeV. The cancellation, however, requires some degree of fine-tuning. Without such cancellations, we would need to require $|\mu|, m_{H^-} > 1$ TeV if we allow deviations from the SM rate to be within 20%.

For $b \to d\gamma$ decay, we should replace $V^*_{ls,ts}$ and λ^2 in Eq. (67) by $V^*_{ld,td}$ and $-\lambda^3 e^{i\phi_1}$, respectively, while Eq. (66) remains unchanged. For $V^*_{ls} V_{mb} \delta^{lm}_{uLL}$ real and negative, it would cancel against charged Higgs contribution for low μ. However, this need not be the case. For example, for $|\mu| = 100$ GeV, $m_{H^-} = 400$ GeV, the cancellation mechanism gives $Br(B \to X_s\gamma)$ within experimental error, while $Br(b \to d\gamma)$ is enhanced by a factor of 2. But in both $b \to s\gamma, d\gamma$ decays, the chargino contribution to asymmetry $a_{M^0\gamma}$ of Eq. (7) is within 2%.

To obtain large asymmetry $a_{\rho\gamma}$, we need a sizable $\sin 2\vartheta$ (see Eq. (8)) which requires $C_{7\gamma}$ and $C'_{7\gamma}$ to be of comparable size. To achieve this the New Physics must have large $C'^{\rm New}_{7\gamma}$ but a relatively small contribution to $C_{7\gamma}$, since the latter already receives a large SM contribution. For example, in case of gluinos, we need large δ^{q3}_{dRR} and/or $\delta^{q3}_{dRL}\, m_{\tilde{g}}/m_b$ and small δ^{q3}_{dLL} and $\delta^{q3}_{dLR}\, m_{\tilde{g}}/m_b$. It is interesting that this indeed can be realized in AHS models. For similar reasons we do not expect a large modification of $C_{7\gamma}$ from squark mixings as was noted in our discussion of charged Higgs effects.

To use the formula of $a_{\rho\gamma}$ we also need to know the phase. The SM phase in $B \to X_d\gamma$ is rather complicated since u and c quark contributions at NLO are not CKM suppressed. However, as shown in [31] the NLO contribution are found to increase the rate by 10 % and the long-distance contribution from intermediate u quarks in the penguin is expected to be small. For exclusive modes, some model estimations (such as from Light-Cone QCD sum rule) give long-distance effect in $B \to \rho\gamma$ and $B \to \omega\gamma$ at about $\mathcal{O}(15\%)$ [67]. Even though long distance physics may enter, it does not enhance $C'^{\rm SM}_{7\gamma}$ [67–69]. For charge B decays the dominant long distance contribution come from weak annihilation diagrams (giving $|C'_{7\gamma}/C^{\rm SM}_{7\gamma}| \sim 4\%$), which is, however, absent in the case of $B^0 \to \rho^0\gamma$ decay.

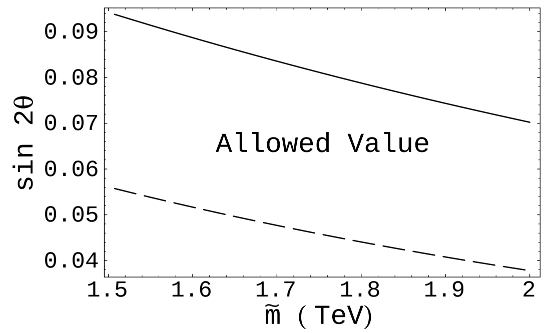

FIG. 9. Asymmetry coefficient $\sin 2\vartheta$ vs. $\widetilde{m}$, for gluino mass $m_{\tilde{g}} = 1.5$ TeV, where solid (dashed) curve correspond to δ^{13}_{RL} and δ^{13}_{RR} having opposite (same) phase.

In Fig. 9, we show gluino and neutralino contributions to the asymmetry coefficient $\sin 2\vartheta$ vs $\widetilde{m}$, with gluino mass $m_{\tilde{g}} = 1.5$ TeV. We use $\mu = 1$ TeV and $\tan\beta = 2$. The asymmetry is generated mainly from the RL mixing induced gluino penguins $\sim 8\%$ and can reach 10% when including RR mixing contributions. The asymmetry can be measured at B Factories, and at future hadron collider B detectors such as LHCb or BTeV. Since $\sin 2\vartheta \sim 2|C'_{7\gamma}/C_{7\gamma}|$ for small ϑ, we obtain $|C'_{7\gamma}/C^{\rm SM}_{7\gamma}|$ of about 4%, for $\widetilde{m} \sim 1.5$ TeV, which is slightly larger than the estimation $\sim 2\%$ for long distance effects from charm penguin [68].

C. Non-standard C-Term and $\tan\beta$ Enhancement

In a previous study, we obtain large or even maximal asymmetry rather easily in $b \to s\gamma$ with sub-TeV superparticle mass scale [30]. Here, the high SUSY scale as required by meson mixings leads to too severe a suppression in $1/G_F\widetilde{m}^2$, as can be seen from Eqs. (58), (59). We find, however, that it is still possible to have large $a_{M^0\gamma}$ when considering *non-standard soft breaking terms* [70]. Non-standard soft breaking terms can survive without inducing quadratic divergence if there is no gauge singlet particles in the low energy spectrum. For scales below the horizontal symmetry breaking scale and the masses of S_i, we are left with particles of the minimal supersymmetry standard model (MSSM). Therefore, by using a low energy effective theory of SUSY, it is legitimate to include these non-standard soft breaking terms, in particular, a non-holomorphic trilinear term, which is called the C-term.

Besides a (standard) A-term, $A_d \langle H_d \rangle Y' \tilde{D}_L \tilde{D}^*_R$, we now allow $(\widetilde{M}^2_d)_{LR}$ to have a non-standard C-term,

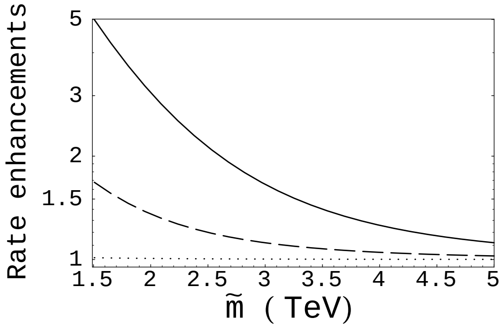

FIG. 10. The rate enhancement of $\mathrm{Br}(B \to X_d\gamma)$ with respect to SM results with non-standard soft breaking terms. The solid, dashed and dotted curves correspond to $\tan\beta = 50$, 20 and 2, respectively. The first curve can enhance the rate up to a factor 5, while the second one can enhance the rate up to a factor 1.8. The enhancement factor for the third one is below 10%.

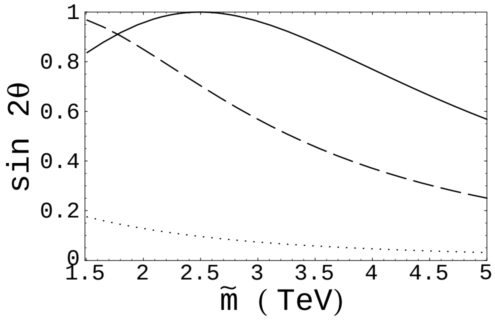

FIG. 11. The CP violating coefficient $\sin 2\vartheta$ with respect to $\widetilde{m}$ for gluino and neutralino contributions. We use $m_{\tilde{g}} = 1.5$ TeV. These curves use the same parameter space as the previous figure. Note that first two curves may reach their maximal value even with a multi-TeV squark mass, while the small $\tan\beta$ case may gives 20 % asymmetry.

$C\langle H_u^*\rangle Y' \tilde{D}_L \tilde{D}_R^*$. It is natural that $A_d \sim C \sim \widetilde{m}$, hence $(\widetilde{M}_d^2)_{LR}^{ij} \sim \widetilde{m}\, M_d^{ij} \tan\beta$. In this way, one gains a $\tan\beta \equiv |\langle H_u^*\rangle/\langle H_d\rangle|$ enhancement factor, while $(\widetilde{M}_u^2)_{LR}^{ij} \sim \widetilde{m}\, M_u^{ij}$ is unaffected. M_q, and hence $U_{qL,R}$, is also unchanged, so the previous result for D^0 mixing remains unchanged. Some zeros in $(\widetilde{M}_q^2)_{LR}$ may also be lifted since these C-terms are no longer holomorphic, but they are still suppressed. We note that the δ_{dRL}^{12} contribution to kaon mixing remain protected by the smallness of $M_d^{21}/\widetilde{m}$. Likewise, for B_d and B_s mixings, $\tan\beta$ enhancement of $\delta_{dLR,RL}^{i3}$ is insufficient to overcome $m_q/\widetilde{m}$ suppression and δ_{dRR}^{i3} still dominates.

We illustrate in Figs. 10 and 11 the ratio $\mathrm{Br}(B \to X_d\gamma)/\mathrm{Br}(B \to X_d\gamma)_{\mathrm{SM}}$ and the coefficient $\sin 2\vartheta$ relevant for mixing dependent CP violation, with respect to the average squark mass $\widetilde{m}$ for $m_{\tilde{g}} = 1.5$ TeV. The solid, dashed and dotted curves correspond to $\tan\beta = 50$, 20 and 2, respectively. The branching ratio can be enhanced by a factor of 5 with respect to the SM value. This can be easily understood by noting that, before introducing the C-term, the RL mixing induced gluino diagrams give $|C'_{7\gamma}/C_{7\gamma}^{\mathrm{SM}}| \sim 4\%$ for $\widetilde{m} \sim 1.5$ TeV. Adding the nonstandard C-term enhances δ_{dRL} by $\tan\beta$ and hence $|C'_{7\gamma}/C_{7\gamma}^{\mathrm{SM}}|$ is brought up to $0.04\tan\beta$. A factor of 5 enhancement in rate follows for $\tan\beta = 50$. Note that $\sin 2\vartheta$ reaches maximum for $\widetilde{m} \sim 2.6$ TeV. The reason is simply because $C'_{7\gamma}$ dominates over $C_{7\gamma}$ for lower $\widetilde{m}$ scale, hence suppresses the asymmetry while enhancing the rate significantly. Since the phase combination $\sin[2\Phi_B - \phi(C_7) - \phi(C'_7)]$ in general should not vanish even if $\phi(C_7^{(\prime)})$ vanishes (because of non-vanishing Φ_B), $a_{\rho^0\gamma}$ could clearly be sizable, which would unequivocally indicate the presence of New Physics.

The CP violating partial rate asymmetry A_{CP} in $b \to d\gamma$ decay is defined as

$$A_{\mathrm{CP}} = \frac{\Gamma(b \to d\gamma) - \bar{\Gamma}(\bar{b} \to \bar{d}\gamma)}{\Gamma(b \to d\gamma) + \bar{\Gamma}(\bar{b} \to \bar{d}\gamma)}$$

$$= \frac{|C_{7\gamma}|^2 + |C'_{7\gamma}|^2 - |\bar{C}_{7\gamma}|^2 - |\bar{C}'_{\gamma 7}|^2}{|C_{7\gamma}|^2 + |C'_{\gamma 7}|^2 + |\bar{C}_{\gamma 7}|^2 + |\bar{C}'_{7\gamma}|^2}, \tag{68}$$

where $\bar{C}_{7\gamma}^{(\prime)}$ are coefficients for $\bar{b}$ decay. To have nonzero A_{CP}, apart from CP phases, one also needs absorptive parts. In the model under consideration, these can come only from the SM contribution with u and c quarks in the loop. The A_{CP} is smaller than the SM one since δ_{LR} which contributes to $C_{7\gamma}$ is much smaller than δ_{RL}, while there is no strong phase in $C'_{7\gamma}$ to contribute to direct CP violation. Therefore New Physics only dilutes the A_{CP} in this case by contributing to the total rate in the denominator of Eq. (68). That is, A_{CP} is reduced by a rate enhancement factor for large $\tan\beta$.

These figures hold also for $b \to s\gamma$ for the other choice of using Eq. (29) with $\tilde{s}$–$\tilde{b}$ but no $\tilde{d}$–$\tilde{b}$ mixing. Allowing for 20% rate uncertainty for the measured $\mathrm{Br}(B \to X_s\gamma)$, we see that for the $\tan\beta = 20$ case, $\widetilde{m} \geq 3$ TeV is allowed, while for heavier squark $\widetilde{m} = 5$ TeV the full range of $2 \lesssim \tan\beta \lesssim 50 \sim m_t/m_b$ is allowed. In these cases $\sin 2\vartheta$ can go up to ~ 0.6. For lighter $\widetilde{m}$ such as 1.5 TeV, the enhancement factor for large $\tan\beta$ starts to break the good agreement between SM and the experimental value of $\mathrm{Br}(B \to X_s\gamma)$, hence it seems one cannot have both large $\tan\beta$ and $\widetilde{m}$, $m_{\tilde{g}}$ too light (approaching TeV).

The interesting case of "strange-beauty" squark [48] where large mixing dependent CP is possible without non-standard C-terms is discussed in next section.

VII. DISCUSSION

We offer to discuss a few miscellaneous items.

A. Re(ε'/ε) and EDM Constraints

Unlike Ref. [34], the AHS model presented here can not be responsible for the large value of $\mathrm{Re}(\varepsilon'/\varepsilon)$. This is because of heavy masses and the suppressed value of $\delta^{12}_{dLR,RL}$, as shown in Table II. However, $\mathrm{Re}(\varepsilon'/\varepsilon)$ does not provide further constraint on δ's, since we already satisfy the most severe case, $\varepsilon_K < 2.268 \times 10^{-3}$. For further discussion on the issue of $\mathrm{Re}(\varepsilon'/\varepsilon)$ in the context of AHS models, see Ref. [71].

There are other horizontal symmetry models that lead to the pattern of quark mass ratios and mixings. In particular, the non-Abelian U(2) horizontal model [72] gives suppressed $\delta^{12}_{dLL,RR}$ and can evade the kaon mixing constraint by a U(2) symmetry with relatively light masses. Mixing angles in fermion mass matrices are in general of the order of the square root of mass ratios [72],

$$M_q = U^\dagger_{qL} M_{\mathrm{diag}} U_{qR}, \tag{69}$$

$$U_{qL,R} = \begin{pmatrix} 1 & s^{12}_{qL,R} & 0 \\ -s^{12}_{qL,R} & 1 & s^{23}_{qL,R} \\ s^{12}_{qL,R} s^{23}_{qL,R} & -s^{23}_{qL,R} & 1 \end{pmatrix}, \tag{70}$$

where

$$s^{23}_{qL} s^{23}_{qR} = \left(\frac{m_2}{m_3}\right)_q, \quad -s^{12}_{qL} = s^{12}_{qR} = \sqrt{\left(\frac{m_1}{m_2}\right)_q}. \tag{71}$$

The model leads to large $\mathrm{Re}(\varepsilon'/\varepsilon)$ [61]. The D-$\bar{D}$ mixing is small in the same way that the kaon mixing constraint is evaded. Mixing in $\tilde{b}$–$\tilde{d}$ is relatively small, but the sparticle scale can be relatively light since kaon and D meson mixing constraints are evaded. This model may also lead to FCNC effects in B system [73].

Let us now consider the constraint from the electric dipole moment (EDM) of the neutron. It is well known that the EDMs of the electron and atoms give severe constraint on SUSY phases [74]. This is a common problem to all SUSY models and is quite independent from FCNC processes considered in this work. The problem should not be worse in our case, and in fact the TeV sparticle scale should loosen the constraint compared with usual considerations.

For the neutron EDM, there are contributions from electric dipole operator, the color dipole operator and the dimension six purely gluonic operator. We expect the first one to be dominant while the others may give comparable contributions and may in some cases loosen the constraint through cancellations [75,76]. In addition, there are two loop contributions [77]. These contributions could become important in the absence of large one loop contributions, such as in SUSY models with massive first and second generation squarks [78]. But since we have potentially large one loop contributions, we use only the electric dipole operator to estimate the order of magnitude bounds. For a more complete recent study of EDM constraint on SUSY models, see Ref. [79].

The neutron EDM can be expressed as $d_n = \frac{1}{3}\eta^E(4d_d - d_u)$, where $\eta^E = 1.53$ is the QCD correction factor [75]. For gluino contribution, we have [41,75],

$$\frac{d_q}{|e|} = \frac{2Q_{\tilde{q}}\alpha_s m_{\tilde{g}}}{3\pi\tilde{m}^2}\, g_4(x_{\tilde{g}\tilde{q}})\, \mathrm{Im}\, \delta^{11}_{qLR}, \tag{72}$$

and for chargino [75],

$$\frac{d_u}{|e|} = \sum_{i=1}^{2} \frac{\alpha_w m_{\tilde{\chi}^+_i}}{4\pi\tilde{m}^2}\, [Q_{\tilde{d}}\, g_4(x_{\tilde{\chi}^+_i \tilde{q}}) + (Q_{\tilde{d}} - Q_u)\, g_3(x_{\tilde{\chi}^+_i \tilde{q}})]\, \mathrm{Im}\, \eta^{\tilde{\chi}^+}_{\tilde{d}}, \tag{73}$$

$$\eta^{\tilde{\chi}^+}_{\tilde{d}} = -\hat{Y}_u \mathcal{V}^*_{j2} (\mathcal{U}^*_{j1} V_{ul} \delta^{lm}_{dLL} V^*_{um} - \mathcal{U}^*_{j2} V_{ul} \delta^{lm}_{dLR} V^*_{um} \hat{Y}_{d_m}). \tag{74}$$

where l, m are summed over three generations of down type squarks. By interchanging $Q_{u,\tilde{d}}$, $\mathcal{U}, V_{ul}\delta^{lm}_{dAB}V^*_{um}$, $\hat{Y}_{d_m}$ with $Q_{d,\tilde{u}}$, $\mathcal{V}, V^*_{u_l d}\delta^{lm}_{uAB}V^*_{u_m d}$, $\hat{Y}_{u_m}$, respectively, one can obtain d_u. Finally, for neutralino contributions, one finds [75],

$$\frac{d_q}{|e|} = \sum_{i=1}^{4} \frac{Q_{\tilde{q}}\alpha_w m_{\tilde{\chi}^0_i}}{4\pi\tilde{m}^2}\, g_4(x_{\tilde{\chi}^0_i \tilde{q}})\, \mathrm{Im}\, \eta^{\tilde{\chi}^0}_{\tilde{q}}, \tag{75}$$

$$\eta^{\tilde{\chi}^0}_{\tilde{q}} = -2G^{*i}_R \delta^{11}_{qRL} G^i_L - \sqrt{2} G^{*i}_R \delta^{11}_{qRR} H^i_R + \sqrt{2} H^{*i}_L \delta^{11}_{dLL} G^i_L + H^{*i}_L \delta^{11}_{qLR} H^i_R. \tag{76}$$

To illustrate the constraint on squark mixing phases we take chargino and neutralino mixing matrices to be real, and discuss the phase of μ later. Requiring d_n to be less than the current experimental bound of $0.63 \times 10^{-25}\, e$ cm [2], we obtain

$$|\mathrm{Im}\, \delta^{11}_{dLR}| \le (2.8,\, 3.4,\, 6.5) \times 10^{-6}, \tag{77}$$

$$|\mathrm{Im}\, \delta^{11}_{uLR}| \le (5.5,\, 6.7,\, 12.6) \times 10^{-6}, \tag{78}$$

for $\widetilde{m} = 1.5$ TeV, $|\mu| = 100$–1000 GeV, and $x_{\tilde{g}\tilde{q}} = (0.3,\, 1,\, 4)$, respectively. These bounds are consistent with Ref. [79]. The bounds come dominantly from gluino contributions and hence insensitive to $\tan\beta$ and $|\mu|$. From chargino contribution alone with above parameter space and $\tan\beta = 50\,(2)$, we have

$$|\mathrm{Im}\, (\sum_{lm} V_{ul} \delta^{lm}_{dLR} V^*_{um} \hat{Y}_{d_m})| \le 0.43 - 0.54\; (0.39 - 0.48),$$

$$|\mathrm{Im}\, (\sum_{lm} V^*_{ld} \delta^{lm}_{uLR} V_{md} \hat{Y}_{u_m})| \le 0.11\; (0.10). \tag{79}$$

The AHS model gives $|\delta^{11}_{dLR}| \sim m_d |A_d(1 + \{\tan\beta\}) - \mu\tan\beta)|/\widetilde{m}^2 \sim |m_d/\widetilde{m}|\, \tan\beta \sim 8.4 \times 10^{-5}(\tan\beta/50)$. Thus, for large $\tan\beta$, we need $\arg(\delta^{11}_{dLR})$ to be less than 0.1 to satisfy the EDM constraint. EDM from Mercury atom gives $d_{H_g} < 2.1 \times 10^{-28}$ e cm [80]. The bounds on $|\mathrm{Im}\delta^{11}_{u,dLR}|$ are one order of magnitude smaller than that from the neutron EDM bounds [79],

$$|\mathrm{Im}\, \delta^{11}_{dLR}| \le (1.1,\, 2.0,\, 4.5) \times 10^{-7}, \tag{80}$$

for 1.5 TeV $\widetilde{m}$. Therefore, $\arg(\delta^{11}_{dLR})$ in this model need to be smaller than 0.01 for large $\tan\beta$.

If μ is complex, it will contribute to $\arg(\delta^{11}_{dLR})$ as $-\arg(\mu)\tan\beta\, m_d|\,\mu|/\widetilde{m}^2$. For large $\tan\beta$ and $|\mu|\sim\widetilde{m}$, we need $\arg(\mu)$ to be less than 0.1 (0.01) from the neutron (Mercury) EDM constraint. One should be more concerned, however, with the presence of $\delta^{11}_{dRR(LL)}$ in Eq. (76). Take δ^{11}_{dRR} for example, we note that it is $\sim\mathcal{O}(1)$ and is not suppressed by quark mass like $\delta^{11}_{dLR(RL)}$. Thus, it will lead to a severe constraint on $\arg(\mu)$. For $m_{\tilde g}=\widetilde{m}=1.5$ TeV, $\tan\beta=2$ and $|\mu|=100$–1000 GeV, we need to have $\arg(\mu)\le 0.03$–0.012. The bound is roughly inversely proportional to $\tan\beta$. For large $\tan\beta$, say 50, we need $\arg(\mu)\le 5\times10^{-4}$ from the neutron EDM constraint. This constraint is quite severe, even for $m_{\tilde g},\widetilde{m}$ at TeV scale. However, the very strong constraint on $\arg(\mu)$ from EDM consideration is a well known problem (see for example, Ref. [79]), and is not aggravated by considerations of FCNC induced by squark mixings, which has been the main focus of our study.

B. Radiative $c\to u\gamma$ and $t\to c\gamma$ Decays

It is of interest to check radiative flavor changing neutral current processes in up type quark decays, since quark-squark alignment has shifted flavor violation to the up-type sector. It is well known that the short distance one-loop $c\to u\gamma$ amplitude is very small in the SM, due to the CKM suppression and the small m_b^2/M_W^2 factor. The amplitude can be raised by 2 orders of magnitude when one considers leading logarithmic QCD corrections involving operator mixings, and further raised by another 3 orders of magnitude when including non-CKM suppressed two-loop diagrams [81]. It is also known that long distance effects are in general large [82].

We can obtain SUSY contribution by using formulas similar to $b\to q\gamma$,

$$\begin{aligned}H_{\rm eff.}=&-\frac{G_F}{\sqrt{2}}\frac{e}{4\pi^2}m_c\,\bar u\left[c_{7\gamma}\,R+c'_{7\gamma}\,L\right]\sigma_{\mu\nu}F^{\mu\nu}c\\&-\frac{G_F}{\sqrt{2}}\frac{g}{4\pi^2}m_c\,\bar u\left[c_{8g}\,R+c'_{8g}L\right]\sigma_{\mu\nu}T^aG_a^{\mu\nu}c,\end{aligned}\qquad(81)$$

Note that here we do not factor out the CKM factor from $c^{\prime}_{7\gamma,8g}$ s (hence use lower case symbol) in the Hamiltonian. For chargino contributions we have,

$$\begin{aligned}c_{7\gamma,\tilde\chi^-}=&\frac{m^2_{\rm W}}{\widetilde{m}^2}\\&\times\Bigg\{\Bigg[\mathcal{U}_{j1}\mathcal{U}^*_{j1}V_{ul}\delta^{lm}_{dLL}V^*_{cm}-\mathcal{U}_{j1}\mathcal{U}^*_{j2}V_{ul}\delta^{lm}_{dLR}V^*_{cm}\hat Y_{d_m}\\&-\mathcal{U}_{j2}\mathcal{U}^*_{j1}V_{ul}\hat Y_{d_l}\delta^{lm}_{dRL}V^*_{cm}+\mathcal{U}_{j2}\mathcal{U}^*_{j2}V_{ul}\hat Y_{d_l}\delta^{lm}_{dRR}V^*_{cm}\hat Y_{d_m}\Bigg]\times\\&\left[Q_dg_2(x_{\tilde\chi^-_j\tilde q})-g_1(x_{\tilde\chi^-_j\tilde q})\right]-\frac{m_{\tilde\chi^-_j}}{m_c}\Big[\mathcal{U}_{j1}\mathcal{V}_{j2}V_{ul}\delta^{lm}_{dLL}\hat Y_cV^*_{cm}\\&-\mathcal{U}_{j2}\mathcal{V}_{j2}V_{ul}\hat Y_{d_l}\delta^{lm}_{dRL}\hat Y_cV^*_{cm}\Big]\left[Q_dg_4(x_{\tilde\chi^-_j\tilde q})-g_3(x_{\tilde\chi^-_j\tilde q})\right]\Bigg\},\end{aligned}$$

$$\begin{aligned}c'_{7\gamma,\tilde\chi^-}=&\hat Y_u\frac{m^2_{\rm W}}{\widetilde{m}^2}\\&\times\Bigg\{\mathcal{V}_{j2}\mathcal{V}^*_{j2}V_{ul}\delta^{lm}_{dLL}\hat Y_cV^*_{cm}\left[Q_dg_2(x_{\tilde\chi^-_j\tilde q})-g_1(x_{\tilde\chi^-_j\tilde q})\right]\\&+\frac{m_{\tilde\chi^-_j}}{m_c}\left[Q_dg_4(x_{\tilde\chi^-_j\tilde q})-g_3(x_{\tilde\chi^-_j\tilde q})\right]\\&\times\left[\mathcal{V}^*_{j2}\mathcal{U}^*_{j2}V_{ul}\delta^{lm}_{dLR}V^*_{cm}\hat Y_{d_m}-\mathcal{V}^*_{j2}\mathcal{U}^*_{j1}V_{ul}\delta^{lm}_{uLL}V^*_{cm}\right]\Bigg\},\end{aligned}\qquad(82)$$

where we sum over l, m for three generations and j for two chargino mass eigenstates. One can obtain $c^{(\prime)}_{8g}$ by replacing $g_{1,3}(x)$ with zero and Q_d by one. Comparing these equations to those in the previous section, we have interchanged $\mathcal{U}$ with $\mathcal{V}$ and modified the charge factor in front of g_i's. In general, $c'_{7\gamma,\chi^-}$ is small due to the smallness of $\hat Y_u$. The chargino loop contribution on $c_{7\gamma}$ is dominated by LL mixing. For $m_{\tilde g}=\widetilde{m}=1.5$ TeV, $|\mu|=100$–1000 GeV, $\tan\beta=2$–50 and $V_{ul}\delta^{lm}_{dLL}V^*_{cm}\sim\lambda$ the chargino loop gives $|c_{7\gamma}|\sim10^{-4}$, which is one order of magnitude below the Cabibbo favored two-loop amplitude. Contributions from LR and RL mixings are smaller by a few orders of magnitude compare to the LL mixing contribution. Therefore, the $\tan\beta$ enhancement effects are small and unable to overcome the heavy superparticle decoupling effects. This is also true in other up-type FCNC processes, such as $t\to c\gamma$ to be discussed later. Charged Higgs contribution is small since we do not have the top quark in the loop.

Formulas for gluino and neutralino contributions are similar to previous section with a trivial modification on neutralino mixing matrices $G_{L,R}$ and $H_{L,R}$. For $\widetilde{m}=m_{\tilde g}=1.5$ TeV, gluino and neutralino loops give $|c_{7\gamma}|\sim10^{-6}$, which contribute to ${\rm Br}(c\to u\gamma)$ at the same order of magnitude of the leading logarithmic SM result.

For the $t\to c\gamma$ case, the SM result is very small, ${\rm Br}(t\to c\gamma)\sim10^{-13}$ [83]. In this model, by using similar formulas for $c^{(\prime)}_{7\gamma}$, we have,

$$\Gamma(t\to c\gamma)=\frac{G_F^2\alpha}{32\pi^4}m_t^5(|c_{7\gamma}|^2+|c'_{7\gamma}|^2).\qquad(83)$$

For the parameter space considered in the previous case, the chargino loop gives ${\rm Br}(t\to c\gamma)\sim10^{-9}$, dominated by chiral enhanced LL mixing. Therefore, it is not sensitive to the non-standard soft breaking term. The rate is still unobservable. The gluino contribution is as small as the SM one. This is in contrast to generic MSSM with non-universal soft squark masses where the rate can be close to experimental bounds [84].

We see that, even though the flavor violation is shifted to the up-type sector, we still do not have large FCNC nor CP violation in t, c decays. This is because of heavy $\widetilde{m}$, $m_{\tilde g}$ masses as required by meson mixings, and absence of enhancement mechanisms in gluino and neutralino diagrams.

C. Light Strange-Beauty Squark?

In this paper we have focused on the general case of naturally large $\tilde{d}_R$-$\tilde{b}_R$ or $\tilde{s}_R$-$\tilde{b}_R$ mixings as a consequence of Abelian flavor symmetry with SUSY. As we have seen, kaon FCNC constraints require texture zeros to remove $\tilde{d}$-$\tilde{s}$ mixing. This is done by QSA, which shifts the source of Cabibbo angle to up-type sector. It then follows that both the K^0 mixing constraint and D^0 mixing bound demand TeV scale SUSY particles. In the case of $\tilde{d}_R$-$\tilde{b}_R$ mixing (mutually exclusive with $\tilde{s}_R$-$\tilde{b}_R$ mixing), B_d mixing constraint also implies TeV scale sparticle masses, and one could get maximal $\sin 2\phi_1$ as suggested by recent Belle result [9].

The $\tilde{s}_R$-$\tilde{b}_R$ mixing case has several special features worthy of note. First, it is maximal, largely because $V_{cb} \sim m_s/m_b \sim \lambda^2$. Second, unlike Δm_{B_d} which is precisely measured already, we only have a lower bound on Δm_{B_s}. In fact, data hints at $\Delta m_{B_s} > \Delta m_{B_s}^{\rm SM}$. From the latter, one cannot draw the conclusion that B_s mixing data demand TeV scale sparticles. From the former, it is intriguing that, in fact, one has a mechanism for the possibility of *one* light squark. As pointed out in Ref. [48], the "democratic" nature of the 2-3 sub-matrix of $\widetilde{M}_{dRR}$ in Eq. (17) not only induces maximal $\tilde{s}_R$–$\tilde{b}_R$ mixing, it could also drive one mass eigenstate $\widetilde{sb}_1$, dubbed the "strange-beauty" [48] squark because it carries both flavors equally, to be much lighter by level splitting. What is rather surprising is that, having $\widetilde{sb}_1$ as light as 100 GeV does not make visible impact on the $b \to s\gamma$ rate. Thus, the light $\widetilde{sb}_1$ scenario survives one of the strongest known constraints on new physics! This has phenomenological bearings.

The general average squark mass scale $\widetilde{m}$ is still fixed by K^0 and D^0 mixings at TeV. But with some tuning in the $\widetilde{M}_{dRR}$ matrix, for example $\widetilde{m}_{23}^2/\widetilde{m}^2 \sim 1$ to λ^3 order, $\widetilde{sb}_1$ can be brought down to 100 GeV. With such large squark mass splittings, the formulas in previous sections do not apply, but one can still follow Ref. [49]. In fact, we find that the $\mathrm{Br}(b \to s\gamma)$ constraint itself can be easily satisfied even if $m_{\widetilde{sb}_1} \to 0$. We note three major consequences of experimental interest: i) Because of low $\widetilde{sb}_1$ mass, sizable C_7' is generated. Although it is subdominant in $b \to s\gamma$ rate, it allows the mixing dependent CP asymmetry, e.g. in $B_s \to \phi\gamma$, to go up to 60%. There is no need to resort to nonstandard C-terms in this case. ii) A light $\widetilde{sb}_1$ squark further enriches B_s mixing and its associated CP phase. Δm_{B_s} could be close to or larger than the SM expectation, and a non-vanishing CP phase in B_s mixing can be measured readily in the moderate x_{B_s} case. iii) A light $\widetilde{sb}_1$ squark clearly offers itself for direct search at future colliders, *in a model where sparticles are otherwise at TeV scale*. In fact, one neutralino, the bino, could also be rather light. A possible decay hence search scenario is $\widetilde{sb}_1 \to (b, s) + \widetilde{\chi}_1^0$ with equal probability of s and b quarks in decay final state. Note that other predictions, such as sizable x_D from SUSY, still hold.

This special scenario, perhaps a bit tuned, seems tailor made for spectacular measurements at the Tevatron and the future LHC. More details and discussions can be found in Ref. [48].

VIII. CONCLUSION

In this work, we make a complete one-loop analysis in SUSY AHS models on FCNC concerning B_d, B_s, K^0, D^0 mixings and $b \to d\gamma$, $s\gamma$ decays. We find that B_d (B_s) and D^0 mixings all receive sizable SUSY contributions even with TeV scale superparticles.

Large off-diagonal elements involving the third generation in the fermion mass matrices follow naturally in AHS models. Hence, flavor mixings involving $\tilde{d}_{jR}$ are naturally prominent. It could be the source for near maximal $\sin 2\phi_1$ given by recent experiments. For $m_{\tilde{q}}$ and $m_{\tilde{g}}$ at TeV scale, the effects could be comparable to SM in B_d (or B_s) mixing, leading to $\phi_{B_d} \neq \phi_1$ (or $\phi_{B_s} \neq 0$), while K^0 mixing and ε_K require quark-squark alignment to make M_d^{12} and M_d^{21} vanish. This shifts V_{us} to u sector, and $\tilde{u}_L$-$\tilde{c}_L$ mixing with masses $\sim$ TeV gives D^0 mixing that is tantalizing close to recent hints from data. With the same squark mixings, the chargino induced contributions to the kaon mixing also points to a TeV scale for superparticle masses, independent of Δm_{B_d} considerations. There is a special variant where a "strange-beauty" squark is driven light by maximal $\tilde{s}_R$-$\tilde{b}_R$ mixing, which can give rise to even more astounding phenomena, but with little impact on B_d system. Otherwise, TeV is in general the preferred sparticle scale in this model.

With such heavy gluino and squarks, one has few other low energy phenomena, and prospects for direct production are depressing. It is possible to have $\sim 10\%$ mixing-dependent asymmetry, $a_{M^0\gamma}$ in $b \to s\gamma$ and $d\gamma$ transitions. In addition, these asymmetries are sensitive to nonstandard soft breaking terms via $\tan\beta$ enhancement, and asymmetries could be up to 60% in $b \to s\gamma$ when the $\mathrm{Br}(B \to X_s\gamma)$ constraint is taken into account. If we insist on non-vanishing M_d^{31}, the kaon mixing constraint requires, indirectly, that s flavor is almost decoupled from the other down flavors. It is then possible that one has SUSY effects in B_d, D^0 but not B_s mixings, while $a_{\rho^0\gamma}$ and $a_{\omega^0\gamma}$ could be maximal with rate enhancements up to a factor of five. Alternatively, effects could concentrate in B_s system and $b \to s\gamma$ (plus D^0 mixing). The phenomenology outlined here can be tested at B factories and the Tevatron in the next few years even if the New Physics scale is so high such that direct searches show no effect.

Acknowledgement. This work is supported in part by the National Science Council of R.O.C. under Grants NSC-89-2112-M-002-063 and NSC-89-2811-M-002-0086, the MOE CosPA Project, and the BCP Topical Program of NCTS. CKC would like to thank the KEK theory group for hospitality.

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